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Production Mechanisms and Signatures of Isosinglet Neutral Heavy Leptons in Z^0 Decays

M. Dittmar *

Department of Physics, University of California-Riverside, USA

A. Santamaria†

Department of Physics, Carnegie Mellon University

Pittsburgh, PA 15213, USA

M. C. Gonzalez-Garcia and J. W. F. Valle‡

Departament de Física Teòrica, Universitat de Valencia

46100 Burjassot, València, SPAIN

Abstract

Neutral Heavy Leptons (NHL) arise in many extensions of the standard electroweak theory such as superstring inspired models. The possibility of gauge singlets NHLS is especially attractive because it gives an explanation to the observed smallness of the neutrino mass. Existing limits on the possible existence of such particles are still fairly poor.

We have investigated isosinglet NHL production and decays within different models. The dominant production cross section is **single production** (*i.e.* $Z^0 \to N + \bar{\nu}$ or $Z^0 \to \bar{N} + \nu$) as a result of mixing with the standard doublet neutrinos. Subsequent NHL decays lead to striking signatures.

Taking into account the expected luminosities and typical detector efficiencies of the different LEP/SLC experiments we conclude that these may discover isosinglet NHLS or else substantially improve and extend present limits on their mass and coupling strength.

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* Also at CERN, 1211 Geneva 23, Switzerland

† Also at Departament de Física Teòrica, Universitat de València and IFIC - CSIC

‡ Also at IFIC, Centre Mixte - CSIC, València, Spain

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1 Introduction

All fundamental electrically *charged* fermions that couple to the $SU(2)_L$ weak interaction have also a *right − handed* component that transforms as a singlet. The simplest way to accomodate the lack of experimental evidence for a non zero neutrino mass is to attribute it to an intrinsic asymmetry in the fermion spectrum as in the standard model *i.e.* to the absence of right handed neutrinos. However there is really no good theoretical basis for this choice.

Small (or zero) neutrino masses [1] can fit naturally in many theoretical contexts that include NHLS such as right handed (RH) neutrinos. These arise in most attempts at unifying the presently observed interactions into a single gauge scheme such as grand unified or superstring inspired models [2].

It was in fact in the previous versions of grand-unification that the existence of right-handed neutrinos combined with the idea of total lepton number violation was suggested as an elegant explanation for the puzzling smallness of neutrino masses, relative to those of the charged fermions [3]. However in most of these models the RH neutrinos are too heavy to be produced in the laboratory.

NHLS also arise in extended electroweak models such as left right symmetry in an attempt to give an explanation to the origin of parity violation in the weak interaction. In these models the RH neutrinos could be accessible to experimental searches in the laboratory and their existence most likely implies non zero neutrino masses.

If there is no sign of left right symmetry at multi-TeV energies one may have the minimal $SU(2) \otimes U(1)$ gauge structure at low energies, but with an extended fermion sector containing the RH neutrino *i.e.* an $SU(2) \otimes U(1)$-based see saw model. In this case the RH neutrinos might be light enough to be produced in the laboratory and again their possible existence is associated to a finite neutrino mass. Consequently the relevant NHL couplings are restricted by the observed smallness of the neutrino masses indicated by laboratory experiments. As a result a possible RH neutrino detection requires a tau neutrino heavier than 1 MeV. Thus in the minimal see saw model there is a contradiction with standard cosmology which requires stable neutrino masses to obey $m_\nu \lesssim 100 eV$ [4]

as well as astrophysics, which puts stringent constraints on possible neutrino decay modes and lifetimes [5]. In ref [5] we analyse possible ways to evade this limitation in models that contain a Majoron [6,7]. The existence of RH neutrinos as required in the see saw model could be associated with the possible observation of neutrino oscillations and flavour violating phenomena such as $\mu \to 3e$ or $\mu \to e\gamma$. However predicted rates for processes such as $\mu \to 3e$ or $\mu \to e\gamma$ are too small to measure in the laboratory, *e.g.* $BR(\mu \to e\gamma) \leq 10^{-15}$ as discussed in ref [12].

Another possibility is to have NHLS such as RH neutrinos *in the absence of neutrino mass effects* [2,8]. In models containing many electroweak fermion singlets neutrino masses can vanish identically thereby avoiding the limitation mentioned above. As a consequence NHL couplings can then be large enough that their production in the laboratory becomes feasible, without conflicting any experimental or cosmological data. In this case there can be no neutrino oscillations *in vacuo* but the non observation of flavour violating phenomena such as $\mu \to 3e$ and $\mu \to e\gamma$ already places important constraints on some of the model parameters [8].

NHLS of any type with mass $M \lesssim 70\ GeV$ might be produced in Z^0 or W decays. Here we concentrate on the spectacular signals that would arise from NHL production and its subsequent decay at the Z^0 peak. The generic features of this process are:

- the NHLS are produced in Intermediate Vector Boson decays due to their admixture in the charged current (CC) and neutral current (NC) of the Standard Model .
- the NHLS decay typically via the CC and NC with rates that scale as M^5 where M is the NHL mass.

In this paper we feel motivated to focus attention on *isosinglet* NHLS rather than a fourth generation sequential *isodoublet* NHL for two main reasons:

- Isosinglet NHLS (in contrast to isodoublets) arise very often in unified and extended electroweak models and their possible existence has an important impact on the physics of neutrino and weak interactions, hopefully serving as guide towards *the* correct theory.

- Isosinglets are *singly*-produced (not *pair* produced) in $e^+ e^-$ colisions, leading to a very characteristic experimental signature.

To set the stage for our discussion we analyse prototype examples of NHL models in section 2.1, followed by a discussion of the general structure of the weak charged and neutral currents present in these models, section 2.2. Laboratory constraints are briefly mentioned in section 2.3. Production and decay mechanisms are studied in section 2.4. The experimental aspects are discussed in section 3.

2 Theory

The possible existence of isosinglet NHLS such as RH neutrinos may be the clue to understand why neutrinos do not acquire mass at the same scale as the charged fermions of the standard model. In a gauge theory *e.g.* the Standard Model, the two issues can not be separated. Other related effects are flavour violation in the lepton sector and/or total lepton number violation (in what follows we will refer to such symmetry as B-L, baryon number B minus total lepton number, $L = L_e + L_\mu + L_\tau$, because this is the combination that often appears in unified models). Many aspects of NHL physics depend on whether this B-L symmetry is broken and in this case, what is the corresponding mass scale and the nature of this symmetry breaking.

In a theory where B-L is a local gauge symmetry we expect it to be broken as there is no evidence for the existence of an intermediate gauge boson lighter than a few TeV that couples to it. This is the situation in left-right theories [10] and some superstring inspired models [11]. This leads to additional interactions of the NHLS with gauge bosons.

When we introduce NHLS in the Standard Model, where B-L is *ungauged*, there is also no reason to keep it as an exact symmetry. If this violation happens in a spontaneous way it leads to a Goldstone boson - the Majoron - which provides neutrinos with new interactions [6,7,12,13].

To simplify the discussion of the different models we first consider here the simplest of all possibilities, where B-L is broken explicitly or unbroken. To keep the theoretical discussion of NHLS completely general it is useful to consider them in the context of the $SU(2) \otimes U(1)$ theory. We consider two classes of models (sections 2.1.1 and 2.1.2) that illustrate the extreme possilities discussed in the introduction. They also illustrate the important role played by B-L symmetry in neutral heavy lepton physics.

2.1 NHL Models

For the sake of simplicity as well as generality we base our discussion on the simplest gauge and Higgs structure of the Standard Model.

Electrically neutral $SU(2)_L$ singlet fermions *e.g.* right-handed neutrino are the only ones that can be added to the Standard Model for which there is an allowed mass term *e.g.* $M_R \nu^c \nu^c$ invariant under both Lorentz symmetry and *standard* gauge symmetry. The Majorana mass leads to total lepton number violation and implies non zero masses for the isodoublet neutrinos.

Another attractive way to introduce isosinglet NHLS into the Standard Model is to keep total lepton number as an exact symmetry responsible for the masslessness of the so far observed neutrinos. This opens a much wider window for lepton physics beyond that possible in the previous case.

The basic Yukawa terms common to all these models are those that lead to charged fermion masses and are completely standard. The models differ only in the neutral fermion sector. We now discuss these models in more detail.

2.1.1 See Saw Model

In this model one uses the following set of (left-handed) fermions (repeated over generation index, i)

$$\begin{pmatrix} \nu_i \\ e_i \end{pmatrix}, \quad {e^c}_i, \quad {\nu^c}_i$$
$$\begin{pmatrix} u_i \\ d_i \end{pmatrix}, \quad u_i^c \quad d_i^c \tag{1}$$

where ν_i^c denotes the RH neutrino of the corresponding generation [§]. The full quark-lepton symmetry in eq. (1) is only apparent since it is not manifest in the actual spectrum of neutrino masses. These are determined from the corresponding neutrino mass matrix (in the basis (ν, ν^c)) [3]

$$\begin{pmatrix} 0 & D \\ D^T & M_R \end{pmatrix} \tag{2}$$

where the matrix D_{ij} is the Dirac mass term for the three right-handed neutrinos, and M_{Rij} is an isosinglet mass, added as a *bare* mass. The resulting neutrino mass after diagonalizing out the heavy fields is

$$M_L{}^{see-saw} = D M_R^{-1} D^T \tag{3}$$

where M_R is the Majorana mass for the right handed neutrinos.

Since ν^c transforms as an electroweak singlet one expects the corresponding coefficient M_R to be large compared with the characteristic scale for the isodoublet mass term $D\nu\nu^c$ which is proportional to the standard Higgs VEV responsible for electroweak symmetry breaking and charged fermion masses. As a result of the coexistence of these two types of mass terms, the physical mass-eigenstate neutrinos are Majorana particles. Their masses are determined by diagonalizing eq. (3) and are suppressed by $O(D/M_R)$ relative to the characteristic mass of the charged fermions. This gives an explanation to the small neutrino masses [3]. The physical mass-eigenstate neutrinos are Majorana particles composed mostly of the isodoublets but containing an admixture of the iso-singlets. The corresponding form of the weak currents is given in section 2.2.1.

From the point of view of the Standard Theory one can formulate a theory of leptons with an arbitrary number of singlets since *e.g.* they do not carry any triangle anomaly [14]. This situation does actually seem to arise in many of the models inspired in the superstring [2,15]. When the number of isosinglets exceeds that of isodoublets new possibilities for lepton physics open up many respects [8] as we will consider in section 3 and 4.

[§] We describe right-handed fermions in charge-conjugate notation.

2.1.2 Total Lepton Number Conserving Model

This model may also be implemented in several gauge models. For simplicity and generality we will discuss it again from the point of view of the standard gauge group $SU(2) \otimes U(1)$. To the fermion content of eq. (1) we now include in addition to right handed neutrinos, an equal number of gauge singlet leptons S_i. In the presence of these extra singlets the masslessness of neutrinos is ensured by imposing the conservation of total lepton number. The same global symmetry keeps neutrinos massless in the standard model. The difference is that here it is imposed while there it is automatic. This restriction leads to the following form form for the neutral mass matrix ¶ (in the basis ν, ν^c, S)

$$\begin{pmatrix} 0 & D & 0 \\ D^T & 0 & M \\ 0 & M^T & 0 \end{pmatrix} \tag{4}$$

It is easy to see that the three light neutrinos are massless *Weyl* neutrinos while the other 6 neutral 2-component leptons combine exactly into 3 heavy *Dirac* fermions. Individual leptonic flavour is however violated ***despite physical neutrinos being strictly massless*** [8,12] as has been discussed in ref [8,12]. The corresponding form of the gauge currents is given in section 2.2.2.

The origin of such extra fermions may be found in some superstring inspired models where often one is led to a mass matrix of the above form [2,15]. The zeroes of these entries arise in some of these models where *e.g.* there are no Higgs fields to provide the usual Majorana mass terms needed in the see-saw mechanism [2,15,18]. We now investigate the consequences of eq. (4) within the framework of the standard $SU(2) \otimes U(1)$ theory. As we will see the phenomenology of NHLS following from eq. (4) is markedly more exciting than that of eq. (2).

¶This form for the mass matrix has been suggested in *different* theoretical models *e.g.* vector-like $SU(2) \otimes U(1)$ models (now excluded by experiment) [16], models incorporating left-right symmetry, where a special discrete symmetry was invoked in order to keep neutrinos massless [17], and in superstring inspired models [15].

2.2 Weak Gauge Currents

Isosinglet neutral heavy leptons (*e.g.* right handed neutrinos) couple to the standard gauge bosons through mixing with ordinary isodoublet neutrinos. This way they can be produced in high energy collisions. The resulting mixing matrix K, describing the charged current leptonic weak interaction in the standard electroweak theory has a *rectangular* form [19]

$$K = (K_L, K_H) \tag{5}$$

In addition, in general there is flavour mixing in the Charged Current of the light neutrinos , even in the very special case where they are kept massless due to the imposition of B-L symmetry [8]. For a n generation model containing m $SU(2) \otimes U(1)$ isosinglets 2-component fermions the matrix K_L is nxn and K_H is nxm. Altough the matrix K can be given explicitly, in general it involves a very large number of parameters: $n(n+2m-1)$ real parameters [19]. This in general far exceeds the number of parameters describing the Charged Current weak interaction of quarks because it includes [19,20]

- light-heavy mixing angles, crucial for our discussion
- possible CP violating phases present in light-heavy mixing
- additional CP violating phases which are present because neutrinos in general have Majorana masses

The corresponding neutral current expressed in terms of mass-eigenstate neutrinos is determined by a related matrix $P = K^\dagger K$ [19] and takes the form

$$P = \begin{pmatrix} K_L^\dagger K_L & K_L^\dagger K_H \\ K_H^\dagger K_L & K_H^\dagger K_H \end{pmatrix} \tag{6}$$

The GIM mechanism is violated due to the admixture of fermions of different weak isospin in the currents. As a result, there are neutral current couplings connecting light to heavy neutrinos. In addition the neutral current couplings involving only light (or heavy) neutrinos among themselves are in general also non-diagonal [19]. The neutral current eq. (6) is determined in terms of the charged current, for which the most general parametrization has been given in ref [19]. Fortunately as we shall see, NHL production on the Z^0 involves very few combinations of these parameters.

We now present a useful form for the charged and neutral currents of each the models considered in sections 2.1.1 and 2.1.2 above. In all cases the strategy is to diagonalize out the heavy fermions using the method of ref [7].

2.2.1 See Saw Model

The charged current in this model can be determined, following ref. [7] as

$$K_L^{seesaw} = U[1 - \frac{1}{2}D^*(M_R^*)^{-1}M_R^{-1}D^T]V_L \tag{7}$$

$$K_H^{seesaw} = UD^*(M_R^*)^{-1}V_H \tag{8}$$

where U, V_L and V_H are arbitrary 3x3 unitary matrices arising from the diagonalization of the various mass matrices. The neutral current in the neutrino sector also becomes non-diagonal when expressed in terms of mass eigenstate neutrinos [19] and can be found from eq. (6). For example the NC couplings of the *light* mass eigenstate doublet neutrinos to the heavy singlet neutrinos is

$$[P_{LH}^{seesaw}]_{ab} \approx [V_L{}^\dagger D^*(M_R^*)^{-1}V_H]_{ab} \tag{9}$$

NHL production from Z^0 decays will proceed through these couplings, see section 2.4. Similarly, we have

$$[P_{LL}^{seesaw}]_{ab} = \delta_{ab} - [V_L{}^\dagger D^*(M_R^*)^{-1}(M_R)^{-1}D^T V_L]_{ab} \tag{10}$$

which shows that the NC in the *light* neutrino sector is in general non-diagonal when expressed in terms of mass eigenstate neutrinos [19]. These couplings determine neutrino decay rates. However these are too small in the present model.

2.2.2 Total Lepton Number Conserving Model

B-L conservation implies that the six 2-component heavy mass-eigenstate leptons following from eq. (4) combine into three *Dirac* particles, so the description of the corresponding mixing matrix K simplifies considerably since it now collapses to a 3×6

matrix. When CP is conserved [||], K can be conveniently written as

$$K_L = UD_A \tag{11}$$

$$K_H = UD_A D_H V_H \tag{12}$$

where the U, V_H are real orthogonal matrices and the D matrices are real, positive and diagonal. They are related by $D_A^2(1 + D_H^2) = 1$ [8]. Since the diagonal matrix D_A is in general non-degenerate, ($D_A \neq 1$) the charged current interactions of the physical mass eigenstate light neutrinos is effectively described by a mixing matrix K_L which is $non - unitary$ [19]. In addition, although the light neutrinos are massless it is not possible to eliminate flavour violating mixing among them. This follows from the non simultaneous diagonalizability of the arbitrary Yukawa couplings originating the D and M entries in eq. (4).

The model defined by the currents in eq. (12) has considerably less parameters than the most general model with the same degrees of freedom as a result of B-L conservation. For example, for the simple case where CP is conserved only 4 parameters in total are needed to describe the weak currents, of a 2 generation model: $\theta_H, \theta_U, h_1, h_2$ arising from V_H, U and D_H respectively.

The corresponding neutral current couplings leading to NHL production can be found from eq. (6),

$$[P_{LH}]_{ab} = [D_A^2 D_H V_H]_{ab} \tag{13}$$

while the corresponding NC coupling of the light neutrino obtained from eq. (6) gives

$$[P_{LL}]_{ab} = [D_A^2]_{ab} \tag{14}$$

Although these couplings are diagonal they differentiate between different light neutrino species.

2.3 Constraints on NHLS

Laboratory experiments constrain isosinglet NHL admixture in the gauge currents, eq. (5). These in turn determine through eq. (18) the NHL couplings strengths relevant for NHL production.

[||]The possibility of CP violation in this system is considered in [9].

First we consider constraints that hold in any of the NHL models discussed above. These come from *low energy* weak decay processes where the neutrinos that can be kinematically produced are only the light ones. NHL admixture in the CC and NC leads in general to violations of universality which limit the attainable values of the K_H matrix elements. For example, from eq. (5) we see that the coupling of a given light neutrino to the corresponding charged lepton is decreased by a certain factor. Constraints on this follow from universality tests *e.g.* β decay-μ decay universality, $e-\mu$ universality in $\pi \to l_2$ decay, τ lifetime, *etc.* Low mass NHL would also be produced in π and K decays, charm and beauty decays, *etc.* and these have been looked for in experiments such as beam dump experiments. These constraints were thoroughly discussed in ref [21] and are summarized below

$$\begin{aligned}(K_H K_H^\dagger)_{ee} &\lesssim 4.3 \times 10^{-2}\\ (K_H K_H^\dagger)_{\mu\mu} &\lesssim 0.8 \times 10^{-2}\\ (K_H K_H^\dagger)_{\tau\tau} &\lesssim 10 \times 10^{-2}\end{aligned} \tag{15}$$

and they can at most reach masses of a few GeV. A tighter constraint extending up to the 10 GeV range has recently been obtained at Fermilab using the wide band neutrino beam [22]. In any case it becomes empty as the the NHL mass approaches 20 GeV.

When the presence of NHLS also engenders non-zero neutrino masses, there are additional, *stronger* limits on the attainable NHL production cross sections. These follow from laboratory and cosmological limits on neutrino masses and lifetimes. Depending on the model these may seriously constrain the discovery potential of NHLS at high energies. This is the case in the see-saw model where the prospects for detecting a NHL signal are rather modest, except when the RH neutrino is acompanied by the existence of a tau neutrino heavier than 1 MeV or so, a value forbidden by cosmology and astrophysics. This may be avoided by introducing a Majoron into the theory [5].

To discover NHLS or to extend limits on their possible existence to high mass values high energy *accelerator* experiments are needed. We are aware of no dedicated search for *isosinglet* NHL, say in $e^+ - e^-$ collider experiments, such as DESY/PETRA or SLAC/PEP nor at hadron colliders, such as the CERN/$S\bar{p}pS$ or FNAL/TEVATRON colliders. As we shall demonstrate Z^0 factories such as LEP will be ideal for NHL searches.

2.4 NHL Production and Decays

First we consider NHL Production Mechanism. From eq. (6) we see that, if at least *one* of the NHLS, say N_a is lighter than the Z^0, the Z^0 will decay into a *single* NHL plus a neutrino with a branching ratio

$$BR(Z \to N_a + \bar{\nu}) = BR(Z \to \nu + \bar{\nu})(\epsilon^\dagger \epsilon)_{aa}(1 - z_a)^2 \, (1 + \frac{1}{2} z_a) \tag{16}$$

where $BR(Z \to \nu + \bar{\nu}) \approx 0.065$ is the branching ratio for producing 1 generation of light isodoublet neutrinos in the standard model. This is plotted in fig 1. Here we have set

$$z_a = \frac{M_a^2}{M_Z^2} \tag{17}$$

where M_a denotes the N_a mass. Here we have redefined the relevant coupling strength of the NHLS to the Z^0 as

$$(\epsilon^\dagger \epsilon)_{aa} = [K_H^\dagger K_H]_{aa}. \tag{18}$$

For each of the above models K_H is given either by eq. (8) or by eq. (12). Here we have summed over all the light neutrino types. There is also a similar decay $Z^0 \to \bar{N}_a + \nu$ with the same branching ratio (we assume CP conservation). In contrast to Z^0 decay to *sequential* doublet NHLS, where *pair* production (if kinematically allowed) would be dominant, here the dominant mode of production is always *single* production. This applies to all iso-singlet models, because the corresponding pair production cross section is suppressed relative to the single production cross section by an additional ϵ^2 factor.

We now come to the expected decay mechanisms of the NHL. For definiteness we consider the case of three generations. Typically isosinglet NHL decay via CC and NC processes, eq. (5) and eq. (6), respectively. The decay rates scale as M^5 where M is the NHL mass. The rate at which new NHL decay channels open up was calculated in detail in [21]. Here we will be mostly concerned with decays of higher mass NHL produced in Z^0 decays and correspondingly we will sum over all 3 flavours, except for the top quark in neutral current processes (assuming $m_t > 46\, GeV$).

Folding the resulting NHL decay branching ratios with the corresponding production rates given from eq. (16) we find the following rates for the combined production and subsequent charged and neutral current mediated decay processes

$$BR(Z^0 \to 2\nu, q, \bar{q}) = 6\mathcal{C}(\epsilon^\dagger\epsilon)_{aa}\left(2(a_u^2+b_u^2)+3(a_d^2+b_d^2)\right)f(z) \tag{19}$$

$$BR(Z^0 \to 4\nu) = 2\mathcal{C}(\epsilon^\dagger\epsilon)_{aa}[\frac{3}{4}f(z)+\frac{1}{4}g(z,z)] \tag{20}$$

$$BR(Z^0 \to \nu, \ell, q, \bar{q}) + BR(Z^0 \to \bar{\nu}, \bar{\ell}, \bar{q}, q) = 6\mathcal{C}(\epsilon^\dagger\epsilon)_{aa}[2f(w,0)+f(w,t)] \tag{21}$$

$$BR(Z^0 \to 2\nu, \bar{\ell}, \ell) = 2\mathcal{C}(\epsilon^\dagger\epsilon)_{aa}\left(3(a_e^2+b_e^2)f(z)+3f(w)-2a_e g(z,w)\right) \tag{22}$$

where the functions $f(z)$, $f(w,t)$ and $g(z,w)$ defined as

$$f(z) = -\frac{2}{z^3}(z^2+3z-6)+\frac{12}{z^4}(1-z)\ln(1-z) \tag{23}$$

$$f(w,t) = \frac{(1-t)}{2w^3}(24-12w(1+t)-2w^2(2-t+2t^2)-15w^3t(1+t)+6t^2w^4)$$
$$-3t^2(4+tw^2)\ln t - \frac{3}{w^4}(1-w)(1-tw)(4-tw^2-t(1+t)w^3-t^2w^4)\ln(\frac{1-tw}{1-w}) \tag{24}$$

$$g(z,w) = \frac{12}{w^3z^3}\left\{-\frac{1}{2}wz(-5z-5w+3zw)-\frac{1}{2}(-3z-2w+rw)z(1-w)\ln(1-w)\right.$$
$$-\frac{1}{2}(-2z-3w+zw)w(1-z)\ln(1-z)+(z+w)(wz-z-w)$$
$$\left.\{\ln(1-z)\ln(\frac{w+z-wz}{z})-Li_2(\frac{w-zw}{w+z-wz})+Li_2(\frac{w}{w+z-wz})\}\right\} \tag{25}$$

take into account propagator and phase space effects in the case of top production through charged currents. In the above equations we have defined the combinations

$$\mathcal{C} \equiv BR(Z \to \nu+\bar{\nu})(1-z)^2(1+\frac{1}{2}z)/\mathcal{N} \tag{26}$$

and

$$\mathcal{N} \equiv 3\{2(a_u^2+b_u^2)+3(a_d^2+b_d^2)\}f(z)+3\{2f(w,0)+f(w,t)\}$$
$$+\frac{1}{4}g(z,z)+\frac{3}{4}f(z)+3(a_e^2+b_e^2)f(z)+3f(w)-2a_e g(z,w) \tag{27}$$

In these equations we have defined z, w and t as

$$z \equiv (M_a/m_Z)^2, \quad w \equiv (M_a/m_W)^2, \quad t \equiv (m_t/M_a)^2 \tag{28}$$

while the a's and b's are given by

$$a_e \equiv -\frac{1}{2}+\sin^2\theta_W, \quad b_e \equiv \sin^2\theta_W,$$
$$a_u \equiv \frac{1}{2}-\frac{2}{3}\sin^2\theta_W, \quad b_u \equiv -\frac{2}{3}\sin^2\theta_W, \quad a_d \equiv -\frac{1}{2}+\frac{1}{3}\sin^2\theta_W, \quad b_d \equiv \frac{1}{3}\sin^2\theta_W \tag{29}$$

In all these processes (except the one mediated by the W and with a top as a final state, if it is light enough) we have neglected final state fermions masses, so there is no phase space suppression. The functions f and g give the effect of the W and Z^0 propagators, *i.e.* the deviations from the pointlike-Fermi interaction. When the NHL mass is small compared with the mass of the W or the Z^0 all these functions are 1, and the decay rate has the same form as in ordinary mu-decay, apart from the existence of additional kinematically allowed decay channels. The same function f occurs both for the processes with charged as well as neutral currents, except that the relevant mass is that of the W in one case and of the Z^0 in the other. This is the reason for the appearance of the two variables z and w. The function g takes into account some interference effects between neutral current diagrams ($g(z,z)$) and between neutral and charged current diagrams ($g(z,w)$) in the decay to leptons. Finally the function $f(w,t)$ is the same as $f(w)$ but including top quark phase space suppression factors in the case of top production in eq. (21) *i.e.* $f(w,0) = f(w)$. For our assumed top quark mass the top contribution becomes negligible, *i.e.* $f(w,t) \ll f(w,0)$. In all of the above we have assumed that the W can not be produced physically *i.e.* $w < 1$ (which implies $z < 1$ as well). The neutral current rates have been calculated in a five quark model, (excluding the top quark).

In eq. (21) and eq. (22) we have summed over all charged lepton flavours. These 3-body decays are common to all NHL models. We have neglected the possible radiative decays $N_a \to \nu + \gamma$ because the corresponding branching ratios are about 4 orders of magnitude smaller than those presented above. Additional NHL decay mechanisms might be present in specific models such as considered in ref [5].

It is necessary experimentally to separate the NHL decays involving $\tau's$ from those involving $\ell = e, \mu$. Thus we now rewrite the corresponding branching ratios for all NHL decay channels with charged leptons in the final state this way. The corresponding branching ratios for the different channels are

$$BR(Z^0 \to \nu, c.lept. \neq \tau, q, \bar{q}) = 6\mathcal{C}\left((\epsilon^\dagger\epsilon)_{aa} - |\epsilon_{\tau a}|^2)(2f(w,0) + f(w,t))\right) \quad (30)$$

$$BR(Z^0 \to \nu, \tau, q, \bar{q}) = BR(Z^0 \to \nu, \bar{\tau}, q, \bar{q}) = 3\mathcal{C}\left(|\epsilon_{\tau a}|^2 (2f(w,0) + f(w,t))\right) \quad (31)$$

$$BR(Z^0 \to 2\nu, 2(c.lept. \neq \tau)) = 4\mathcal{C}\{(\epsilon^\dagger\epsilon)_{aa}(a_e^2+b_e^2)f(z) + ((\epsilon^\dagger\epsilon)_{aa} - |\epsilon_{\tau a}|^2)f(w)$$
$$-a_e((\epsilon^\dagger\epsilon)_{aa} - |\epsilon_{\tau a}|^2)g(z,w)\} \quad (32)$$

$$BR(Z^0 \to 2\nu, \tau, c.lept \neq \tau) = BR(Z^0 \to 2\nu, \bar{\tau}, c.lept \neq \tau) = \mathcal{C}\{((\epsilon^\dagger\epsilon)_{aa} + |\epsilon_{\tau a}|^2)f(w)\} \quad (33)$$

$$BR(Z^0 \to 2\nu, \tau, \bar{\tau}) = 2\mathcal{C}\{(\epsilon^\dagger\epsilon)_{aa}(a_e^2+b_e^2)f(z) + |\epsilon_{\tau a}|^2 f(w) - 2a_e |\epsilon_{\tau a}|^2 g(z,w)\} \quad (34)$$

We will now discuss in detail the expected signatures and discovery capabilities of NHLS at LEP/SLC, taking into account explicitly the expected detection efficiencies for the various channels. It will be clear from next section that NHL discovery through leptonic decays will be the most promising.

3 Single NHL Signature at the Z^0 Peak

We now discuss the expected experimental signatures of a heavy isosinglet neutral lepton at LEP and SLC. These machines will run, after some initial searches for new particles and a precise scan to determine the Z^0 mass and width at the Z^0 peak. One expects on average 2-3 million of Z^0 decays per year and experiment. One expects to collect a total of about 10 million Z^0 decays before the LEP energy can be increased to $\sqrt{s} = 190\,GeV$. This high event rate makes it feasible to look for rare Z^0 decays like the possible production of an iso-singlet NHL N with its subsequent decay. As discussed above, such NHL is expected to decay via the charged or neutral weak current

$$N_a \to \text{"}Z\text{"}\nu\,;\ \text{"}Z\text{"} \to ee,\ \mu\mu\ etc. \quad (35)$$

$$N_a \to \text{"}W\text{"}e,\mu,\tau\,;\ \text{"}W\text{"} \to e\nu,\ qq\ etc. \quad (36)$$

The total leptonic branching ratio is roughly 30 per cent. The experimental signature will therefore be very unbalanced events with a huge amount of missing transverse momentum. The pure leptonic decays will have low mass pairs of leptons with high momentum and small opening angle. The decays into quarks will result into jets, which for a mass $M \approx 10-20GeV$ will result in collimated jets, "mono-jets" which have relatively small masses and large missing momenta. Similar signatures have been discussed for all kinds of reactions involving the production and detection of supersymmetry [23]. Very similar signatures and cross sections have been investigated for reactions like $e^+e^- \to \tilde{Z}\tilde{\gamma}$ and the subsequent $\tilde{Z}$ decays into leptons or quarks and an invisible photino.

It has been found that the detection of these objects, even though they are quite rare is relatively easy for the leptonic decay modes and more complicated for the jet like final states.

To investigate the possible sensitivity of the LEP/SLC experiments we have simulated the decay $Z^0 \to N + \bar{\nu}$ and the subsequent different decay modes of the N. For its decay and the possible subsequent jet fragmentation we have used the LUND-Monte Carlo frame [24].

The angular distribution with respect to the beam direction for spin 1/2 particles is $1 + cos^2\theta$. The easiest signatures are the ones where the N decays into pure lepton final states. To determine the possible experimental sensitivity to this reaction we have used some easy acceptance criteria which can be fulfilled, according to their technical proposals, by any LEP/SLC experiments, see for example ref [25]. We have used the following criteria:

1. We have required events with two oppositely charged particles and no isolated high energy γ in the final state. At least one of these particles has to be an electron or a μ.

2. Each of these charged particles should be well inside the acceptance of the tracking detector, $-0.9 < \cos\theta < 0.9$ and has to have a momentum above $5\,GeV$. In the events where one of the decay products might be a τ which will decay further, we require a minimum momentum of $1\,GeV$ for the final state charged particles.

3. We required that the missing transverse momentum of these two particles exceeds 5 GeV.

4. Finally, we require that these two tracks are acoplanar in the plane transverse to the beam, their opening angle should be smaller than 120 degrees.

These criteria are not optimised for all possible masses but are chosen to reject the standard leptonic decays of the Z^0. Especially the acoplanarity requirement rejects very efficiently the background from $\tau^+\tau^-$ events. These cuts provide background free event signals over a mass range for the N between $10\,GeV$ and $70\,GeV$. The efficiency for accepting this decay mode of the N varies between 50-30 percent for N masses between $10\,GeV$ and $70\,GeV$ for the decay $N \to \tau + e(\mu) + \nu$. Using the same criteria for the N decays without $\tau's$ we find roughly 10-20 percent higher efficiencies.

For an observation of this rare decay we required to find 10 events. Unfolding the efficiency and the corresponding branching ratios we can determine the sensitivity for this rare decay mode as a function of the observed number of Z^0 decays, the mass M of this heavy N and its mixing parameter ϵ. Our results are shown in figures 2, 3, 4, and 5.

We have investigated also the decays into hadronic final states which are expected to have about a factor of 3 higher branching ratio. However, to reject efficiently the background from hadronic events, we find that an almost complete coverage of the detectors with electromagnetic and hadron calorimeters is needed. A good resolution for total jet energies or at least the possibility of rejecting jets which contain high energy neutrons or $K_L's$ is essential if background suppression factors of 10^5 have to be reached. Since the total energy measurement and neutron detection is very detector specific we did not investigate this further. We want to point out however, that this decay mode will increase the sensitivity to this possible rare decay. Also, the additional observation of these decay modes together with the mixed leptonic final states will prove that the new signature cannot be explained with supersymmetry or similar other exotic objects. Finally, the observation of a few 100 events with the decay $N \to qqe(\mu)$ would allow to measure the mass of the N and by looking at the forward backward charge asymmetry of the positive (negative) charged lepton one could measure the vector coupling of the $N's$ **

**For our study we have ignored the forward backward asymmetry which is expected to be ≈ 0.15 for ν events

4 Discussion

LEP/SLC experiments can probe isosinglet NHL mass and coupling strength parameters far beyond the range accesible to other laboratory experiments.

As we have discussed, the existence of NHLS is usually related to the small masses of the isodoublet neutrinos. These are tightly constrained by laboratory as well as standard big bang cosmology. In fact from fig 3 we see that in the see-saw model a measurable NHL signal at LEP/SLC is inconsistent with cosmology and astrophysics since a tau neutrino heavier than an MeV or so is required. For this reason we have not considered this model in any further detail †† and have focused our attention on the model of section 2.1.2.

In the model we have discussed in section 2.1.2 the imposition of total lepton number conservation avoids the limitations found in the see saw scheme since neutrino remain massless to the extent that this symmetry holds exactly.

Thus in the large class of models suggested here the existence of NHL opens up an interesting window for a rich spectrum of physics beyond the Standard Model including a broad range of possible new phenomena such as:

- universality violation in low energy weak decays [21,8], potentially larger in the model of section 2.1.2.
- unstable neutrino decaying invisibly via the neutral current (as $\nu' \to 3\nu$) or via the CC and NC as $\nu' \to e^+e^-\nu$ [19,12] possible in the model of section 2.1.1
- effects of lepton flavour violation related to
  - non zero neutrino masses *e.g.* neutrino oscillations in the see saw model of section 2.1.1
  - the effective non-orthogonality of the flavour neutrinos following from NHL admixture in the charged and neutral currents as in the model of section 2.1.2 [8]

††This limitation of the see saw scheme may be avoided by including a Majoron. A detailed discussion of this question is given in ref [5]

- leptonic CP violation [9]
- non-standard neutrino propagation properties [26] *etc.*

Some of these effects could be observed in the next generation of experiments. For example lepton flavour violating effects could be much larger than expected in most $SU(2) \otimes U(1)$-based NHL models and are not directly related to the neutrino mass.

These include *measurable* branching ratios for low energy lepton flavour violating processes such as $\mu \to e + \gamma$, $\mu \to 3e$, *etc.*

In addition flavour violation can occur in the decays of the Z^0 [8] causing non-standard Z^0 decays *e.g.* $Z^0 \to e\bar{\tau}$, and $Z^0 \to \mu\bar{\tau}$ (plus their conjugates). These processes can be enhanced without unnatural fine-tuning of the parameters nor conflict with any limits on neutrino mass. A detailed analysis of their observability at LEP/SLC is now in progress [27].

5 Conclusions

Iso-singlet NHLS in the mass range between $10\,GeV$ and $70\,GeV$ could be discovered at LEP, through their leptonic decays. If no signal is seen one will be able, for the first time to set limits on their admixture in the weak currents over this mass range. The limits depend on the luminosity to be reached and less sensitively, on the mass of the N. Our results are shown in figure 3 for the case of the see saw model and figures 4 and 5 for the case of the total lepton number conserving model. In the first case NHL production may only be observed for natural parameter choices if the τ neutrino has a mass larger than 1 MeV. However this is inconsistent with astrophysical and cosmological limits on ν_τ mass and lifetimes [4]. In the second case measurable production rates are fully consistent with all known facts about neutrino physics.

NHLS could also manifest their existence through indirect effects as discussed above.

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Figure Captions

Fig. 1 : Branching ratio for $Z^0 \to N + \bar{\nu}$ plus $Z^0 \to \bar{N} + \nu$ divided by the coupling strength parameter $(\epsilon^\dagger \epsilon)_{aa} = [K_H^\dagger K_H]_{aa}$ as a function of the NHL mass.

Fig. 2 : NHL coupling strength parameter $(\epsilon^\dagger \epsilon)_{aa} = [K_H^\dagger K_H]_{aa}$ that would lead to 10 NHL decay events plotted as a function of the NHL mass for different assumed luminosities: (a) 10^6 Z's, (b) 10^5 Z's and (c) 10^4 Z's. Here we have included all visible NHL decays.

Fig. 3 : Total visible NHL decay branching ratio in the see-saw model for a given value of the corresponding light neutrino mass. The solid curve is for a NHL mass $M = 10\,GeV$ and the dashed one is for $M = 80\,GeV$. Measurable rates are incompatible with the cosmological critical density constraint on light neutrino masses and lifetimes. The branching ratio becomes very small for low neutrino masses because the NHL coupling ϵ becomes too small.

Fig. 4 : Detection efficiencies as a function of the NHL mass. Curve (a) is for the decay into $\tau\bar{\tau}\nu$, (b) is for decay into $\ell\bar{\tau}\nu$ plus $\bar{\ell}\tau\bar{\nu}$ where $\ell = e, \mu$, and (c) is for decay into $\ell\bar{\ell}\nu$ plus $\bar{\ell}\ell\bar{\nu}$.

Fig. 5 : Limits on NHL coupling strength parameter that can be reached for different number of Z^0's plotted as a function of the NHL mass. This is calculated for the τ type NHL neglecting mixing between different lepton families. The analysis for the decay signatures for NHLS associated with the e or μ family (and including mixing) can be obtained from eq. (32), eq. (33) and eq. (34). Here we have included only leptonic final states, including NHL decays into τ's with the corresponding detection efficiencies. We have not indicated the present constraints from experiments at low energy.

Figure 1

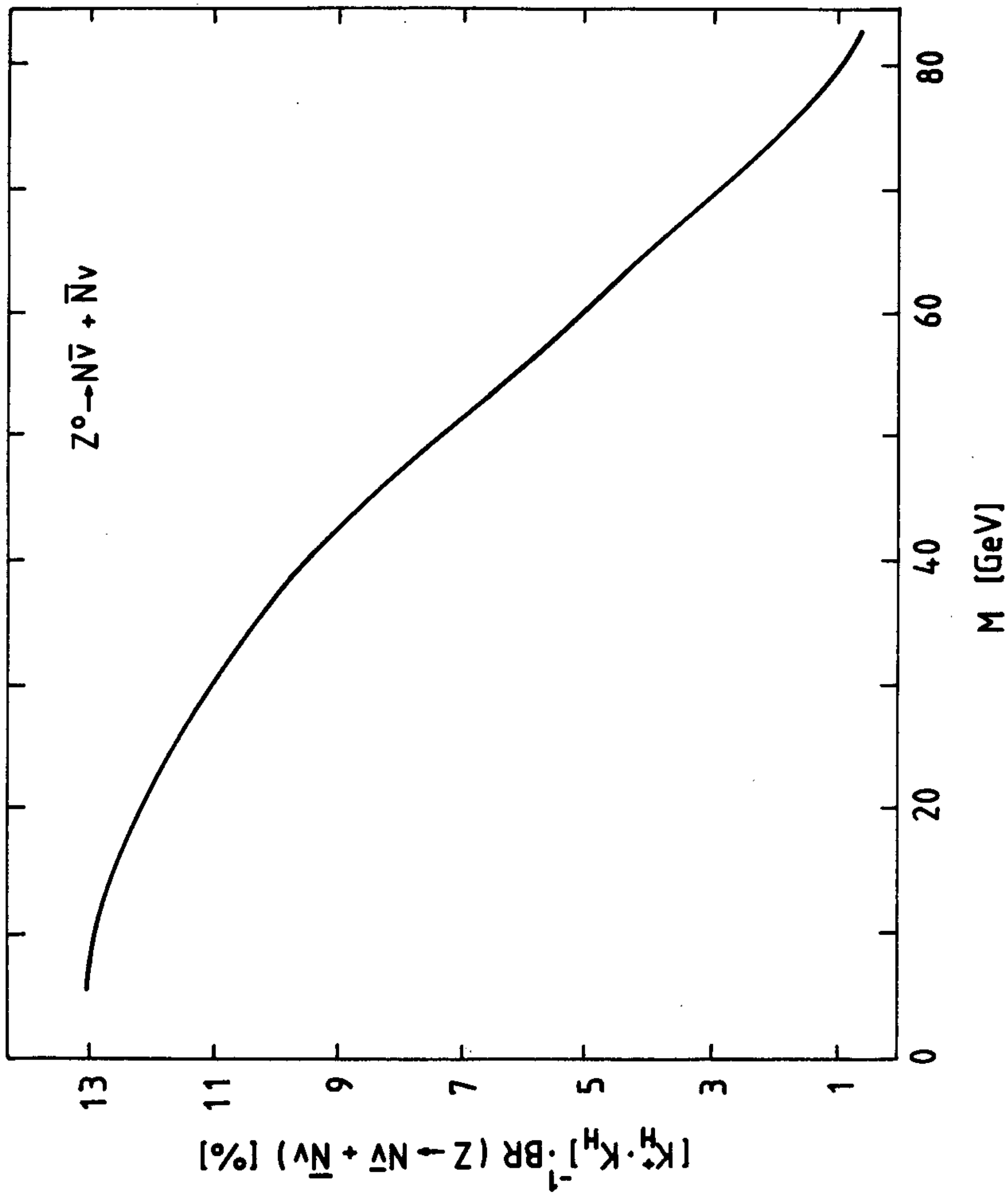
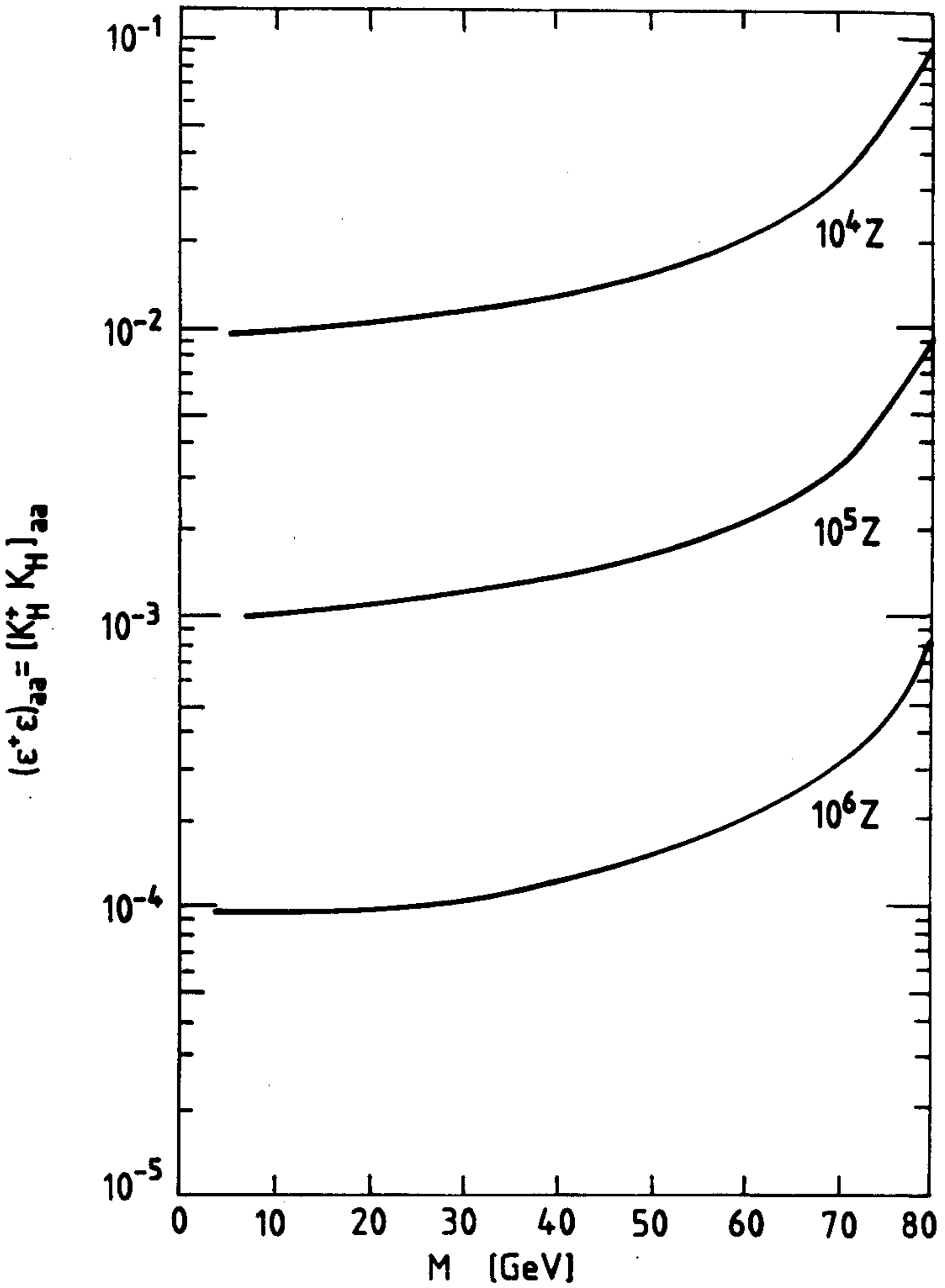


Figure 2



10⁻¹
10⁻²
10⁻³
10⁻⁴
10⁻⁵
(ε⁺ε)aa = [K⁺H KH]aa
10⁴Z
10⁵Z
10⁶Z
0
10
20
30
40
50
60
70
80
M [GeV]

Figure 3

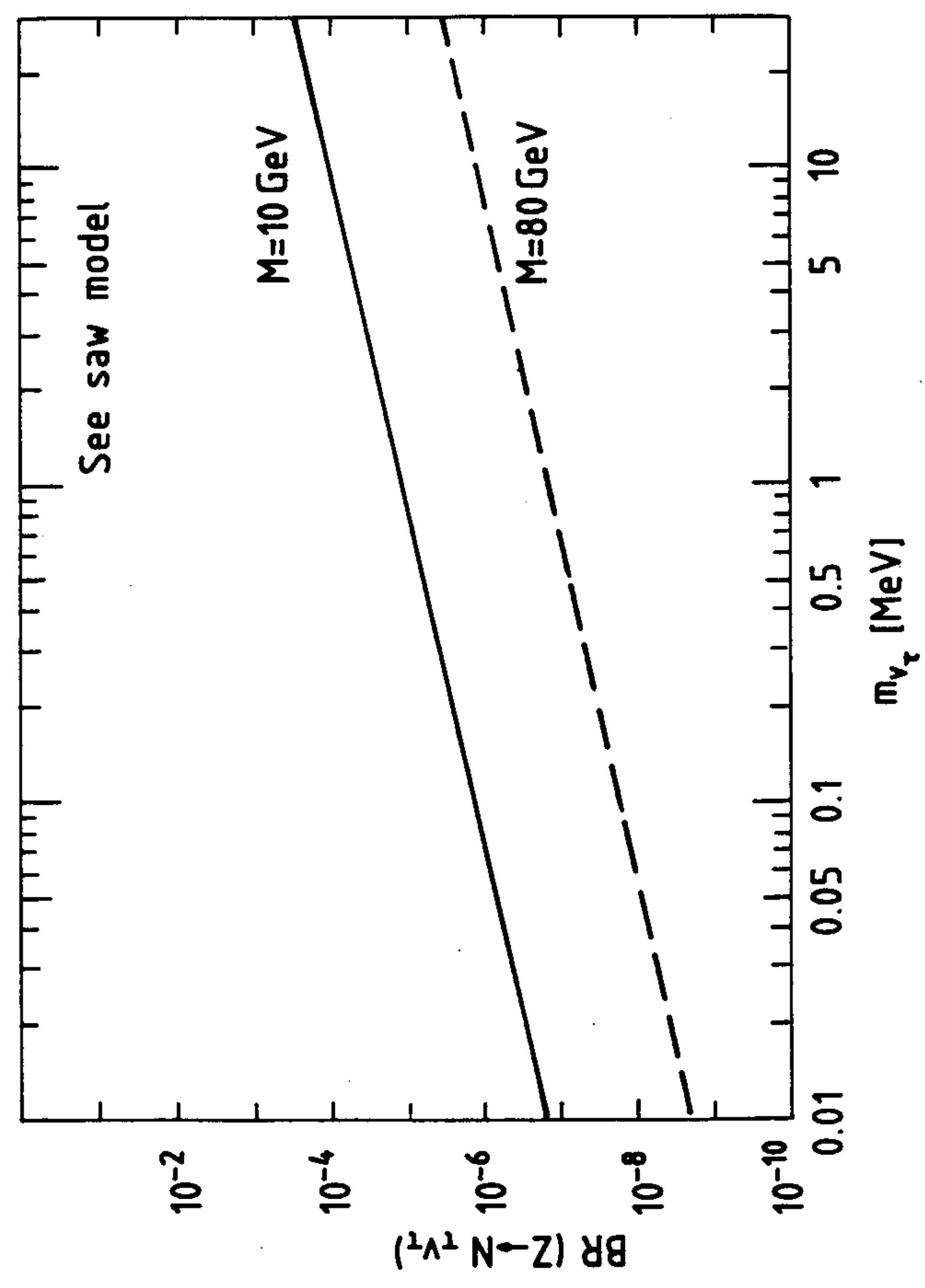

See saw model
M=10 GeV
M=80 GeV
BR (Z→N τν_τ)
10^-2
10^-4
10^-6
10^-8
10^-10
0.01
0.05
0.1
0.5
1
5
10
m_ντ [MeV]

Figure 4

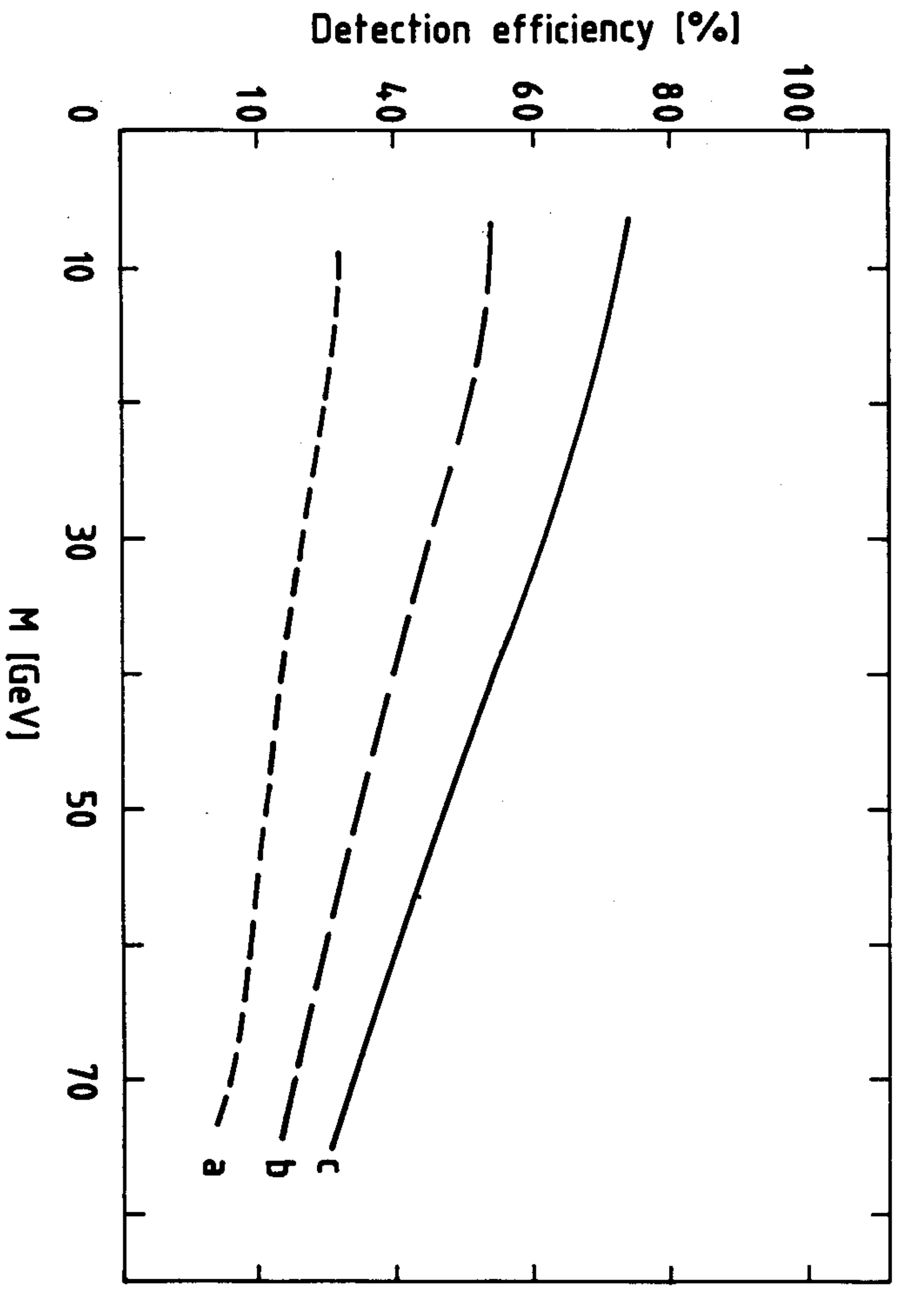
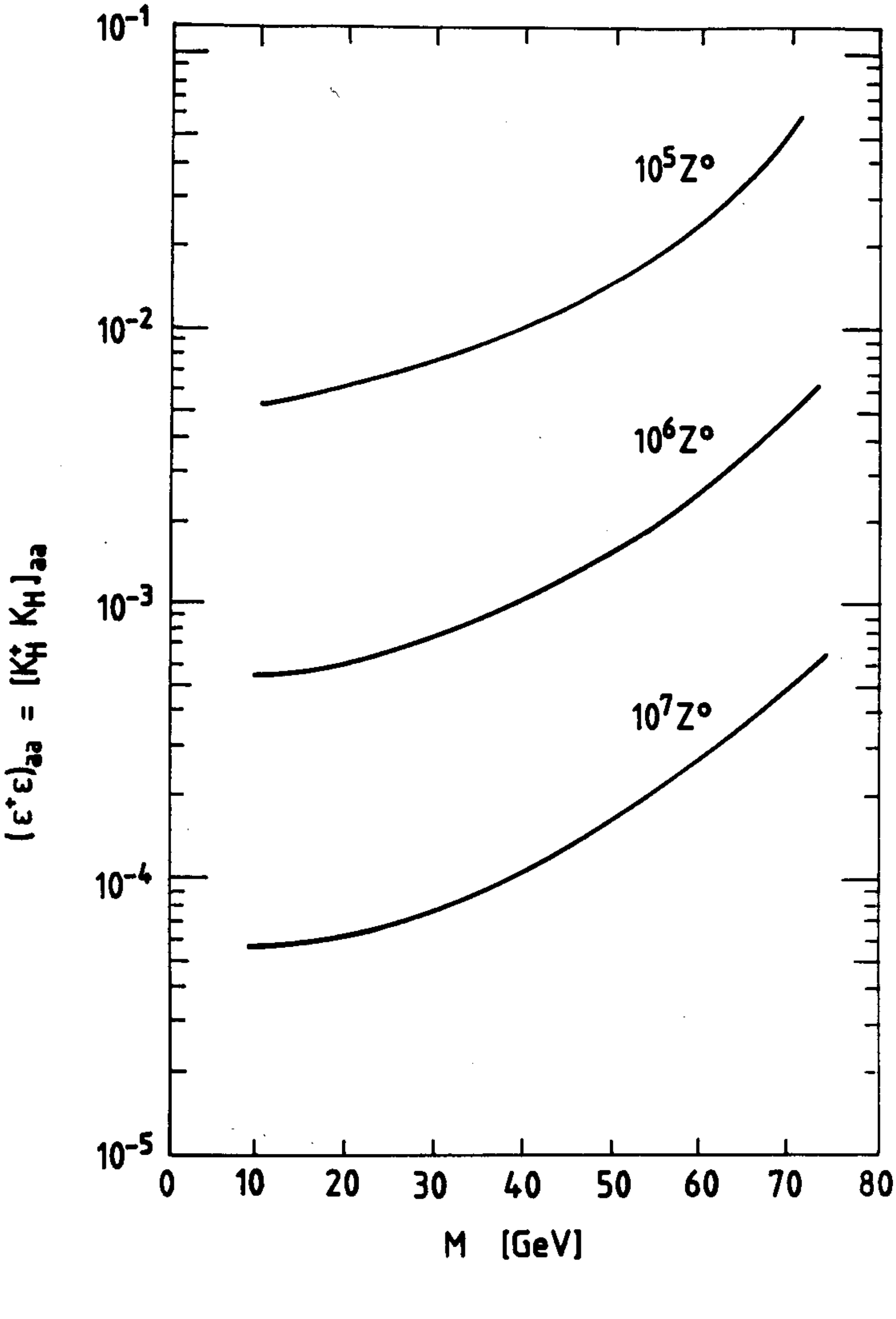


Figure 5


10^{-1}
10^{-2}
10^{-3}
10^{-4}
10^{-5}
0
10
20
30
40
50
60
70
80
M [GeV]
$(\varepsilon^{+}\varepsilon)_{aa} = [K_H^{+} K_H]_{aa}$
$10^5 Z^0$
$10^6 Z^0$
$10^7 Z^0$