Physics Beyond the Desert

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1. Introduction

Neutrinos are the only apparently massless electrically neutral fermions in the SM and the only ones without $SU(2) \otimes U(1)$ singlet partners. It is rather mysterious why neutrinos seem to be so special when compared with the other fundamental fermions. Indeed one of the most unpleasant features of the SM is that the masslessness of neutrinos is not dictated by an underlying principle, such as that of gauge invariance in the case of the photon: the SM simply postulates that neutrinos are massless and, as a result, many of their properties are trivial. If massive, neutrinos would present another puzzle, of why are their masses so much smaller than those of the charged fermions. The fact that neutrinos are the only electrically neutral elementary fermions may hold the key to the answer, namely neutrinos could be Majorana fermions, the most fundamental ones. In this case the suppression of their mass could be associated to lepton number conservation, as actually happens in many extensions of the SM.
Although attractive, the seesaw mechanism [1] is by no means the only way to generate neutrino masses. There are many other attractive possibilities, some of which do not require any large mass scale. The extra particles required to generate the neutrino masses have masses accessible to present experiments [2].

It is also quite plausible that B-L or lepton number, instead of being part of the gauge symmetry [1, 3] may be a spontaneously broken global symmetry. The scale at which such a symmetry gets broken does not need to be high, as in the original proposal [1], but can be rather low, close to the weak scale [3]. Such a low scale for lepton number breaking could have important implications not only in astrophysics and cosmology but also in particle physics.

This large diversity of possible schemes and the lack of a theory for the Yukawa couplings imply that present theory is not capable of predicting the scale of neutrino masses any better than it can fix the masses of the other fermions, like that of the muon. As a result one should at this point turn to experiment.

From the observational point of view non-zero neutrino masses now seem required in order to account for the data on solar and atmospheric neutrinos, as well as cosmological data on the amplitude of primordial density fluctuations which suggest the need for hot dark matter in the universe. I briefly overview the present observational limits and hints in favour of massive neutrinos, and make a few general remarks about the theoretical models.

Turning to the electroweak breaking sector, I mention some of the physics motivations and potential of various extensions of the SM, with emphasis on supersymmetry (SUSY) with broken R–parity and majoron extensions of the SM. Some of the related phenomena are deeply related to the neutrino sector of the theory. I discuss the some aspects of the physics of invisibly decaying Higgs bosons and its impact on Higgs boson searches at accelerators, such as LEP II. Second, I discuss a few aspects of models beyond the Minimal Supersymmetric Standard Model (MSSM) phenomenology, in which R parity is violated, as well as some of the associated laboratory signatures.

Many of the new signatures may be accessible to experiments performed at accelerators or at underground installations, thus illustrating the complementarity between these two approaches in the search for signals beyond the SM.

2. Models of Neutrino Mass

One of the most attractive approaches to generate neutrino masses is from unification. Indeed, in trying to understand the origin of parity violation in the weak interaction by ascribing it to a spontaneous breaking phenomenon, in the same way as the W and Z acquire their masses in the SM, one arrives at the so-called left-right symmetric extensions such as $SU(2)_L \otimes SU(2)_R \otimes U(1)_B$, $SU(4) \otimes SU(2) \otimes SU(2)$ [4] or $SO(10)$ [5], in some of which the masses of the light neutrinos are obtained by diagonalizing the following mass matrix in the basis $\nu, \nu^c$

$$
\begin{bmatrix}
0 & D \\
D^T & M_R
\end{bmatrix}
$$

(1)
where $D$ is the standard $SU(2) \otimes U(1)$ breaking Dirac mass term and $M_R = M^T_R$ is the isosinglet Majorana mass. In the seesaw approximation, one finds

$$M_{\text{eff}} = -DM_R^{-1}DT.$$  \hspace{1cm} (2)

In general, however, this matrix also contains a $\nu\nu$ term \cite{7} whose size is expected to be also suppressed by the left-right breaking scale. As a result one is able to explain naturally the relative smallness of neutrino masses. Even though it is natural to expect $M_R$ to be large, its magnitude heavily depends on the model. Moreover the $M_R$ may have different possible structures in flavour space (so-called textures) \cite{8}. As a result one can not make any real prediction for the corresponding light neutrino masses that are generated through the seesaw mechanism. In fact this freedom has been exploited in model building in order to account for an almost degenerate neutrino mass spectrum \cite{4}.

Although very attractive, unification is not the only way to generate neutrino masses. There are many other attractive schemes which do not require any large mass scale. For example, it is possible to start from an extension of the lepton sector of the $SU(2) \otimes U(1)$ theory by adding a set of two 2-component isosinglet neutral fermions, denoted $\nu^c_i$ and $S_i$, to each generation $i$. In this case there is an exact L symmetry that keeps neutrinos strictly massless, as in the SM. The conservation of total lepton number leads to the following form for the neutral mass matrix (in the basis $\nu, \nu^c, S$)

$$\begin{pmatrix}
0 & D & 0 \\
D^T & 0 & M \\
0 & M^T & 0
\end{pmatrix}$$ \hspace{1cm} (3)

This form has also been suggested in various theoretical models \cite{10}, including superstring inspired models. In the latter case the zeros of eq. (3) naturally arise due to the absence of Higgs fields to provide the usual Majorana mass terms, needed in the seesaw scheme \cite{11}. Clearly, one can easily introduce non-zero masses in this model through a $\mu SS$ term that could be proportional to the vacuum expectation value (VEV) of a singlet field $\sigma$ \cite{12}. In contrast to the seesaw scheme, the neutrino masses are directly proportional to $\langle \sigma \rangle$. This model provides a conceptually simple and phenomenologically rich extension of the SM, which opens the possibility of many new phenomena. These have to do with neutrino mixing, universality, flavour and CP violation in the lepton sector \cite{13,14}, as well as direct effects associated with Neutral Heavy Lepton (NHL) production at high energy colliders \cite{15}. A remarkable feature of this model is the possibility of non-trivial neutrino mixing despite the fact that neutrinos are strictly massless. This tree-level effect leads to a new type of resonant neutrino conversion mechanism that could play an important role in supernovae \cite{16,17}. Moreover, there are loop-induced lepton flavour and CP non-conservation effects whose rates are precisely calculable \cite{13,14,18}. I repeat that this is remarkable due to the fact that physical light neutrinos are massless, as in the SM. This feature is the same as what happens in the supersymmetric mechanism of flavour violation \cite{19}. Indeed, in the simplest

3
case of SU(5) supergravity unification, there are flavour violating processes, like $\mu \to e\gamma$, despite the fact that in SU(5) neutrinos are protected by B-L and remain massless. The SUSY mechanism and that of eq. (3) differ in that the lepton flavour violating processes are induced in one case by NHL loops, while in SUSY they are induced by scalar boson loops. In both cases the particles in the loops have masses at the weak scale, leading to branching ratios $[13, 14, 18, 20, 21]$ that are sizeable enough to be of experimental interest $[22, 23, 24]$.

There is also a large variety of radiative schemes to generate neutrino masses. The prototype models of this type are the Zee model and the model suggested by Babu $[2]$. In these models lepton number is explicitly broken, but it is easy to realize them with spontaneous breaking of lepton number. For example in the version suggested in ref. $[25]$ the neutrino mass arises from the diagram shown in Fig. (1).

The seesaw and the radiative mechanisms of neutrino mass generation may be combined. Supersymmetry with broken R-parity also provides a very elegant mechanism for the origin of neutrino mass, as well as mixings $[26]$. Here I focus on the simplest unified supergravity version of the model with bilinear breaking of R–parity, characterized universal boundary conditions for the soft breaking parameters $[26, 27]$. In this model the tau neutrino $\nu_\tau$ acquires a mass, due to the mixing between neutrinos and neutralinos given in the matrix

$$
\begin{bmatrix}
M_1 & 0 & -\frac{1}{2}g'v_1 & \frac{1}{2}g'v_2 & -\frac{1}{2}g'v_3 \\
0 & M_2 & \frac{1}{2}g'v_1 & -\frac{1}{2}g'v_2 & \frac{1}{2}g'v_3 \\
-\frac{1}{2}g'v_1 & \frac{1}{2}g'v_1 & 0 & -\mu & 0 \\
\frac{1}{2}g'v_2 & -\frac{1}{2}g'v_2 & -\mu & 0 & \epsilon_3 \\
-\frac{1}{2}g'v_3 & \frac{1}{2}g'v_3 & 0 & \epsilon_3 & 0
\end{bmatrix}
$$

(4)

This model contains only one extra free parameter in addition to those of the minimal supergravity model, as the $\epsilon_3$ and the $\nu_3$ are related by a minimization condition. Contrary to a popular misconception, the bilinear violation of R–parity implied by the parameter $\epsilon_3$ is physical and can not be rotated away $[28]$. In fact, what happens in this model with universal conditions for the soft breaking parameters is that the value of $\epsilon_3$ is induced radiatively,
due to the effect of the non-zero bottom quark Yukawa coupling \( h_b \) in the running of the renormalization group equations from the unification scale down to the weak scale $^{[28]}$. This makes $\epsilon_3$ and $m_{\nu_\tau}$ calculable. Thus eq. (4) is analogous to a see-saw type matrix eq. (1), in which the $M_R$ lies at the weak scale (neutralinos), while the rôle of the Dirac entry $D$ is played by the $\epsilon_3$, which is, in a sense, a radiatively induced quantity. From this point of view, the mechanism is a hybrid see-saw like scheme, with naturally suppressed Majorana $\nu_\tau$ mass induced by the mixing between weak eigenstate neutrinos and Higgsinos or gauginos. The $\nu_\tau$ mass induced this way depends quadratically on an effective parameter $\xi$ defined as $\xi \equiv (\epsilon_3 v_1 + \mu v_3)^2$ characterizing the violation of either through $v_3$ or $\epsilon_3$. In Fig. (2) we display the allowed values of $m_{\nu_\tau}$ which clearly can be quite low, due to the possible cancellation between the two terms in $\xi$. In unified supergravity models with universal soft masses this cancellation happens automatically and is, as mentioned, calculable in terms of $h_b$. Notice that $\nu_e$ and $\nu_\mu$ remain massless in this approximation. They get masses either from scalar loop contributions in Fig. (3) $^{[29, 28]}$ or by mixing with singlets in models with spontaneous breaking of R-parity $^{[31]}$. It is important to notice that even when $m_{\nu_\tau}$ is small, many of the
corresponding R-parity violating effects can be sizeable. An obvious example is the fact that the lightest neutralino decay will typically decay inside the detector, unlike the case of the MSSM.

Other than the seesaw scheme, none of the above models requires a large mass scale. In all of them one can implement the spontaneous violation of the global lepton number symmetry leading to neutrino masses that scale directly proportional to the lepton-number scale or some positive power of it, in contrast to the original majoron model [4]. Such low-scale models are very attractive and lead to a richer phenomenology, as the extra particles required have masses at scales that could be accessible to present experiments. One remarkable example is the possibility invisibly decaying Higgs bosons [3].

The above discussion should suffice to illustrate the enormous freedom and wealth of phenomenological possibilities in the neutrino sector. These reach well beyond the realm of conventional neutrino experiments, including also signatures that can be probed, though indirectly, at high energy accelerators. An optimist would regard as very exciting the fact that the neutrino sector may hold so many experimental possibilities, while a pessimist would be discouraged by the fact that one does not know the relevant scale responsible for neutrino mass, nor the underlying mechanism. Last but not least, one lacks a theory for the Yukawa couplings. As a consequence neutrino masses are not predicted and it is up to observation to search for any possible clue. Given the theoretical uncertainties in predicting neutrino masses from first principles, one must turn to observation. Here the information comes from laboratory, astrophysics and cosmology.

2.1. Laboratory Limits

The best limits on the neutrino masses can be summarized as [31]:

\[ m_{\nu_e} < \sim 5 \text{ eV}, \quad m_{\nu_\mu} < \sim 170 \text{ keV}, \quad m_{\nu_\tau} < \sim 18 \text{ MeV} \]  

These are the most model-independent of the laboratory limits on neutrino mass, as they follow purely from kinematics. The limit on the \( \nu_e \) mass comes from beta decay, that on the \( \nu_\mu \) mass comes from PSI (90 % C.L.) [32], with further improvement limited by the uncertainty in the \( \pi^- \) mass. On the other hand, the best \( \nu_\tau \) mass limit now comes from high energy LEP experiments [33] and may be substantially improved at a future tau-charm factory [34]. In connection with tritium beta decay limit [35] even though the negative \( m^2 \) value has now been clarified, there are still un-understood features in the spectrum, probably of instrumental origin. Further results from the Mainz experiment are awaited.

Additional limits on neutrino masses follow from the non-observation of neutrino oscillations. The most stringent bounds come from reactor oscillation experiments [36] (\( \bar{\nu}_e - \nu_x \) oscillations). Here we highlight the recent results of the first long-baseline reactor neutrino oscillation experiment Chooz [37]. There are also stringent bounds from meson factory oscillation experiments (KARMEN [38], LSND [39]) and from high-energy accelerator experiments E531 and E776 [40] (\( \nu_\mu - \nu_\tau \)). A search for \( \nu_\mu \) to \( \nu_e \) oscillations has now been
reported by the LSND collaboration using $\nu_\mu$ from $\pi^+$ decay in flight [41]. An excess in the number of beam-related events from the $C(\nu_e, e^-)X$ inclusive reaction is observed. The excess cannot be explained by normal $\nu_e$ contamination in the beam at a confidence level greater than 99%. If interpreted as an oscillation signal, the observed oscillation probability of $(2.6 \pm 1.0 \pm 0.5) \times 10^{-3}$ is consistent with the previously reported $\bar{\nu}_\mu$ to $\bar{\nu}_e$ oscillation evidence from LSND. Another recent result comes from NOMAD and rules out part of the LSND region. The future lies in searches for oscillations using accelerator beams directed to far-out underground detectors, with very good prospects for the long-baseline experiments proposed at KEK, CERN and Fermilab.

If neutrinos are of Majorana type a new form of nuclear double beta decay would take place in which no neutrinos are emitted in the final state, i.e. the process by which an $(A, Z-2)$ nucleus decays to $(A, Z)+2 e^-$. In such process one would have a virtual exchange of Majorana neutrinos. Unlike ordinary double beta decay, the neutrino-less process violates lepton number and its existence would indicate the Majorana nature of neutrinos. Because of the phase space advantage, this process is a very sensitive tool to probe into the nature of neutrinos.

Present data place an important limit on a weighted average neutrino mass parameter $\langle m \rangle \lesssim 1 - 2$ eV. The present experimental situation as well as future prospects is illustrated in Fig. (4), taken from ref. [42]. Note that this bound depends to some extent on the relevant nuclear matrix elements characterising this process [43]. The parameter $\langle m \rangle$ involves both neutrino masses and mixings. Thus, although rather stringent, this limit may allow very large neutrino masses, as there may be strong cancellations between different neutrino types. This may happen automatically in the presence of suitable symmetries. For example, the decay vanishes if the intermediate neutrinos are Dirac-type, as a result of the corresponding lepton number symmetry [44].
Neutrino-less double beta decay has a great conceptual importance. It has been shown \cite{13} that in a gauge theory of the weak interactions a non-vanishing $\beta\beta_0^\nu$ decay rate requires neutrinos to be Majorana particles, \emph{irrespective of which mechanism} induces it. This is important since in a gauge theory neutrino-less double beta decay may be induced in other ways, e.g. via scalar boson exchange.

2.2. Limits from Cosmology

There are a variety of cosmological arguments that give information on neutrino parameters. In what follows I briefly consider the critical density and the primordial Nucleosynthesis arguments.

2.2.1. The Cosmological Density Limit

The oldest cosmological bound on neutrino masses follows from avoiding the overabundance of relic neutrinos \cite{16}

$$\sum m_{\nu_i} \lesssim 92 \Omega_\nu h^2 \text{eV}, \quad (6)$$

where $\Omega_\nu h^2 \leq 1$ and the sum runs over all species of isodoublet neutrinos with mass less than $O(1 \text{ MeV})$. Here $\Omega_\nu = \rho_\nu / \rho_c$, where $\rho_\nu$ is the neutrino contribution to the total density and $\rho_c$ is the critical density. The factor $h^2$ measures the uncertainty in the present value of the Hubble parameter, $0.4 \leq h \leq 1$, and $\Omega_\nu h^2$ is smaller than 1. For the $\nu_\mu$ and $\nu_\tau$ this bound is much more stringent than the laboratory limits eq. (5).

Apart from the experimental interest \cite{34}, an MeV tau neutrino also seems interesting from the point of view of structure formation \cite{17}. Moreover, it is theoretically viable as the constraint in eq. (6) holds only if neutrinos are stable on the relevant cosmological time scale. In models with spontaneous violation of total lepton number \cite{4} there are new interactions of neutrinos with the majorons which may cause neutrinos to decay into a lighter neutrino plus a majoron, for example \cite{18},

$$\nu_\tau \to \nu_\mu + J \quad (7)$$

or have sizeable annihilations to these majorons,

$$\nu_\tau + \nu_\tau \to J + J \quad (8).$$

The possible existence of fast decay and/or annihilation channels could eliminate relic neutrinos and therefore allow them to have higher masses, as long as the lifetime is short enough to allow for an adequate red-shift of the heavy neutrino decay products. These 2-body decays can be much faster than the visible modes, such as radiative decays of the type $\nu' \to \nu + \gamma$. Moreover, the majoron decays are almost unconstrained by astrophysics and cosmology (for a detailed discussion see ref. \cite{16}).

A general method to determine the majoron emission decay rates of neutrinos was first given in ref. \cite{19}. The resulting decay rates are rather model-dependent and will not be discussed here. Explicit neutrino decay lifetime estimates are given in ref. \cite{30, 18, 50}. The conclusion is that there are many ways to make neutrinos sufficiently short-lived and that all mass values consistent with laboratory experiments are cosmologically acceptable.
2.2.2. The Nucleosynthesis Limit  There are stronger limits on neutrino lifetimes or annihilation cross sections arising from cosmological nucleosynthesis. Recent data on the primordial deuterium abundance [51] have stimulated a lot of work on the subject [52]. If a massive $\nu_\tau$ is stable on the nucleosynthesis time scale, ($\nu_\tau$ lifetime longer than $\sim 100$ sec), it can lead to an excessive amount of primordial helium due to their large contribution to the total energy density. This bound can be expressed through an effective number of massless neutrino species ($N_{\nu}$). If $N_\nu < 3.4 - 3.6$, one can rule out $\nu_\tau$ masses above 0.5 MeV [53, 54]. If we take $N_{\nu} < 4$ the $m_{\nu_\tau}$ limit loosens accordingly. However it has recently been argued that non-equilibrium effects from the light neutrinos arising from the annihilations of the heavy $\nu_\tau$ ’s make the constraint a bit stronger in the large $m_{\nu_\tau}$ region [55]. In practice, all $\nu_\tau$ masses on the few MeV range are ruled out. One can show, however that in the presence of $\nu_\tau$ annihilations the nucleosynthesis $m_{\nu_\tau}$ bound is substantially weakened or eliminated [56].

Fig. 4 gives the effective number of massless neutrinos equivalent to the contribution of a massive $\nu_\tau$ majoron model with different values of the coupling $g$ between $\nu_\tau$’s and $J$’s, expressed in units of $10^{-5}$. For comparison, the dashed line corresponds to the SM $g = 0$ case. One sees that for a fixed $N_{\nu}^{max}$, a wide range of tau neutrino masses is allowed for large enough values of $g$. No $\nu_\tau$ masses below the LEP limit can be ruled out, as long as $g$ exceeds a few times $10^{-4}$. One can express the above results in the $m_{\nu_\tau} - g$ plane, as shown in figure 6. One sees that the constraints on the mass of a Majorana $\nu_\tau$ from primordial nucleosynthesis can be substantially relaxed if annihilations $\nu_\tau\bar{\nu_\tau} \leftrightarrow JJ$ are present. Moreover the required values of $g(m_{\nu_\tau})$ are reasonable in many majoron models [48, 56, 57]. Similar depletion in massive $\nu_\tau$ relic abundance also happens if the $\nu_\tau$ is unstable on the nucleosynthesis time scale [58] as will happen in many majoron models.

Figure 5. A heavy $\nu_\tau$ annihilating to majorons can lower the equivalent massless-neutrino number in nucleosynthesis.
2.3. Limits from Astrophysics

There are a variety of limits on neutrino parameters that follow from astrophysics, e.g. from the SN1987A observations, as well as from supernova theory, including supernova dynamics \[59\] and from nucleosynthesis in supernovae \[60\]. Here I briefly discuss three recent examples of how supernova physics constrains neutrino parameters.

It has been noted a long time ago that, in some circumstances, massless neutrinos may be mixed in the leptonic charged current \[16\]. Conventional neutrino oscillation searches in vacuo are insensitive to this mixing. However, such neutrinos may resonantly convert in the dense medium of a supernova \[16, 17\]. The observation of the energy spectrum of the SN1987A $\bar{\nu}_e$'s \[61\] may be used to provide very stringent constraints on massless neutrino mixing angles, as seen in Fig. (7). The regions to the right of the solid curves are forbidden, those to the left are allowed. Massless neutrino mixing may also have important implications for \(r\)-process nucleosynthesis in the supernova \[60\]. For details see ref. \[17\].

Another illustration of how supernova restricts neutrino properties has been recently considered in ref. \[62\]. There flavour changing neutral current (FCNC) neutrino interactions were considered. These may induce resonant massless-neutrino conversions in a dense supernova medium, both in the massless and massive case. The restrictions that follow from the observed $\bar{\nu}_e$ energy spectra from SN1987A and the supernova $r$-process nucleosynthesis provide constraints on supersymmetric models with \(R\) parity violation, which are much more stringent than those obtained from the laboratory. In Fig. (8) and Fig. (9) we display the constraints on explicit \(R\)-parity-violating FCNCs in the presence of non-zero neutrino masses in the hot dark matter eV range. As seen from Fig. (8) and Fig. (9) they isfavour a leptoquark interpretation of the recent HERA anomaly.

As a final example of how astrophysics can constrain neutrino properties we consider the case of resonant $\nu_e \to \nu_s$ and $\bar{\nu}_e \to \bar{\nu}_s$ conversions in supernovae, where $\nu_s$ is a sterile neutrino \[63\], which we assume to be in the hot dark matter mass range. The implications of such a
Figure 7. SN1987A bounds on massless neutrino mixing.

scenario for the supernova shock re-heating, the detected $\bar{\nu}_e$ signal from SN1987A and for the $r$-process nucleosynthesis hypothesis have been recently analysed [63]. In Fig. (10), taken from [63], we illustrate the resulting constraints on mixing and mass difference for the $\nu_e - \nu_s$ system that follow from the supernova shock re-heating argument. Notice that for the case of $r$-process nucleosynthesis there is an allowed region for which the $r$-process nucleosynthesis can be enhanced.

3. Indications for Neutrino Mass

So far most of positive hints in favour of nonzero neutrino rest masses come from astrophysics and cosmology, with a varying degree of theoretical assumptions. We now turn to these.

3.1. Dark Matter

Considerations based on structure formation in the Universe have become a popular way to argue in favour of the need of a massive neutrino [64]. Indeed, by combining the observations of cosmic background temperature anisotropies on large scales performed by the COBE satellite [65] with cluster-cluster correlation data e.g. from IRAS [60] one finds that it is
Supernovae and FCNC neutrino interactions. 

\begin{figure}[h]
\centering
\includegraphics[width=0.5\textwidth]{figure8}
\caption{Supernovae and FCNC neutrino interactions.}
\end{figure}

not possible to fit well the data on all scales within the framework of the simplest cold dark matter (CDM) model. The simplest way to obtain a good fit is to postulate that there is a mixture of cold and hot components, consisting of about 80\% CDM with about 20\% hot dark matter (HDM) and a small amount in baryons. The best candidate for the hot dark matter component is a massive neutrino of about 5 eV. It has been argued that this could be the tau neutrino, in which case one might expect the existence of $\nu_e \rightarrow \nu_\tau$ or $\nu_\mu \rightarrow \nu_\tau$ oscillations. Searches are now underway at CERN \cite{67}, with a similar proposal at Fermilab. This mass scale is also consistent with the hints in favour of neutrino oscillations reported by the LSND experiment \cite{39}.

3.2. Solar Neutrinos

The averaged data collected by the chlorine, Kamiokande, as well as by the low-energy data on pp neutrinos from the GALLEX and SAGE experiments still pose a persisting puzzle, now re-confirmed by the first 200 days of Super-Kamiokande (SK) data \cite{68}. The most recent data can be summarised in Fig. \cite{11} where the theoretical predictions refer to the BP95 SSM prediction of ref. \cite{69}. For the gallium result we have taken the average of the GALLEX and the SAGE measurements.

The totality of the data strongly suggests that the solar neutrino problem is real, that the simplest astrophysical solutions are ruled out, and therefore that new physics is needed \cite{70}. The most attractive possibility is to assume the existence of neutrino conversions involving very small neutrino masses. In the framework of the MSW effect \cite{71} the required solar neutrino parameters $\Delta m^2$ and $\sin^2 2\theta$ are determined through a $\chi^2$ fit of the experimental
Supernovae and FCNC neutrino interactions.

In Fig. (12), taken from ref. [72], we show the allowed two-flavour regions obtained in an updated MSW analysis of the solar neutrino data including the recent SK 200 days data, in the BP95 model for the case of active neutrino conversions. The analysis of spectral distortion as well as day-night effect plays an important role in ruling out large region of parameters. Compared with previously, the impact of the recent SK data is felt mostly in the large mixing solution which, however, does not give as good a fit as the small mixing solution, due mostly to the larger reduction of the $^7$Be flux found in the later. The most popular alternative solutions to the solar neutrino anomaly include the MSW sterile neutrino conversions, as well as the just-so or vacuum oscillation solution. Recent fits have also been given including the recent SK data [72].

A theoretical point of direct phenomenological interest for Borexino is the study of the possible effect of random fluctuations in the solar matter density [73]. The existence of noise fluctuations at a few percent level is not excluded by the SSM nor by present helioseismology studies. They may strongly affect the $^7$Be neutrino component of the solar neutrino spectrum so that the Borexino experiment should provide an ideal test, if sufficiently small errors can be achieved. The potential of Borexino in "testing" the level of solar matter density fluctuations is discussed quantitatively in ref. [74].

3.3. Atmospheric Neutrinos

Water Cerenkov underground experiments, Kamiokande, Superkamiokande and IMB show a clear deficit in the expected flux of atmospheric $\nu_\mu$’s relative to that of $\nu_e$’s that would be produced from conventional decays of $\pi$’s, $K$’s as well as secondary muon decays [68, 75]. This
is also seen in the iron calorimeter Soudan2 experiment. Although the predicted absolute fluxes of neutrinos produced by cosmic-ray interactions in the atmosphere are uncertain at the 20% level, their ratios are expected to be accurate to within 5%. Other experiments, such as Frejus and NUSEX, have not found a firm evidence, but they have much larger errors. It is tempting therefore to take seriously the evidence for an atmospheric neutrino deficit and to ascribe to neutrino oscillations. In ref. [76] the impact of recent experimental results on atmospheric neutrinos from experiments such as Superkamiokande and Soudan on the determinations of atmospheric neutrino oscillation parameters is considered for the $\nu_\mu \rightarrow \nu_\tau$ channel. In performing this re-analysis theoretical improvements in flux calculations as well as neutrino-nucleon cross sections have been taken into account. The relevant allowed regions of oscillation parameters have been determined from a fit of the various data. One of the new features that arises from the inclusion of the Superkamiokande data is that the best fit value of the $\Delta m^2$ is lower than previously obtained. In fact for the $\nu_\mu \rightarrow \nu_e$ channel the allowed region is almost totally ruled out by the recent Chooz data [37].

4. Reconciling Present Hints

4.1. Almost Degenerate Neutrinos

The above observations from cosmology and astrophysics do seem to suggest a theoretical puzzle. As can easily be understood just on the basis of numerology, it seems rather difficult to reconcile the three observations discussed above in a framework containing just the three known neutrinos. The only possibility to fit these observations in a world with just the three known neutrinos is if all of them have nearly the same mass $\sim 2$ eV [77]. This can
be arranged, for example in general seesaw models which also contain an effective triplet VEV [3, 4] contributing to the light neutrino masses. This term should be added to eq. (2). Thus one can construct extended seesaw models where the main contribution to the light neutrino masses ($\sim 2$ eV) is universal, due to a suitable horizontal symmetry, while the splittings between $\nu_e$ and $\nu_\mu$ explain the solar neutrino deficit and that between $\nu_\mu$ and $\nu_\tau$ explain the atmospheric neutrino anomaly [9].

4.2. Four-Neutrino Models

A simpler alternative way to fit all the data is to add a fourth neutrino species which, from the LEP data on the invisible Z width, we know must be of the sterile type, call it $\nu_s$. The first scheme of this type gives mass to only one of the three neutrinos at the tree level, keeping the other two massless [78].

Two basic schemes of this type that keep the sterile neutrino light due to a special symmetry have been suggested. In addition to the sterile neutrino $\nu_s$, they invoke additional Higgs bosons beyond that of the SM, in order to generate radiatively the scales required for
Figure 12. Solar neutrino parameters for active MSW conversions.

![Solar neutrino parameters for active MSW conversions.](image)

Figure 13. "Heavy" Sterile 4-Neutrino Model

the solar and atmospheric neutrino conversions. In these models the $\nu_s$ either lies at the dark matter scale as illustrated in Fig. (14) or, alternatively, at the solar neutrino scale. In the first case the atmospheric neutrino puzzle is explained by $\nu_\mu$ to $\nu_s$ oscillations, while in the second it is explained by $\nu_\mu$ to $\nu_\tau$ oscillations. Correspondingly, the deficit of solar neutrinos is explained in the first case by $\nu_e$ to $\nu_\tau$ oscillations, while in the second it is explained by $\nu_e$ to $\nu_s$ oscillations. In both cases it is possible to fit all observations together. However, in the first case there is a clash with the bounds from big-bang nucleosynthesis. In the latter case the $\nu_s$ is at the MSW scale so that nucleosynthesis limits are satisfied. They nicely agree with the best fit points of the atmospheric neutrino parameters from Kamiokande. Moreover, it can naturally fit the hints of neutrino oscillations of the LSND experiment. Another theoretical possibility is that all active neutrinos are very light, while the sterile neutrino $\nu_s$ is the single neutrino responsible for the dark matter.
4.3. MeV Tau Neutrino

An MeV range tau neutrino is an interesting possibility to consider for two reasons. First, such mass is within the range of the detectability, for example at a tau-charm factory [34]. On the other hand, if such neutrino decays before the matter dominance epoch, its decay products would add energy to the radiation, thereby delaying the time at which the matter and radiation contributions to the energy density of the universe become equal. Such delay would allow one to reduce the density fluctuations on the smaller scales purely within the standard cold dark matter scenario, and could thus reconcile the large scale fluctuations observed by COBE [65] with the observations such as those of IRAS [66] on the fluctuations on smaller scales.

In ref. [82] a model was presented where an unstable MeV Majorana tau neutrino naturally reconciles the cosmological observations of large and small-scale density fluctuations with the cold dark matter model (CDM) and, simultaneously, with the data on solar and atmospheric neutrinos discussed above. The solar neutrino deficit is explained through long wavelength, so-called just-so oscillations involving conversions of $\nu_e$ into both $\nu_\mu$ and a sterile species $\nu_s$, while the atmospheric neutrino data are explained through $\nu_\mu \rightarrow \nu_e$ conversions. Future long baseline neutrino oscillation experiments, as well as some reactor experiments will test this hypothesis. The model assumes the spontaneous violation of a global lepton number symmetry at the weak scale. The breaking of this symmetry generates the cosmologically required decay of the $\nu_\tau$ with lifetime $\tau_{\nu_\tau} \sim 10^2 - 10^4$ seconds, as well as the masses and oscillations of the three light neutrinos $\nu_e$, $\nu_\mu$, and $\nu_s$ required in order to account for the solar and atmospheric neutrino data. One can verify that the big-bang nucleosynthesis constraints [53, 54] can be satisfied in this model.

5. Electroweak Symmetry Breaking

A basic ingredient in the SM is the breaking of the electroweak symmetry via the Higgs mechanism. If indeed the Higgs boson exists as an elementary particle, apart from its direct
search, the main task in these investigations is the study of supersymmetric extensions of the SM and the corresponding experimental searches at high energy accelerators.

5.1. Supersymmetry and the MSSM

The main theoretical motivation for SUSY are that it allows for a stable hierarchy between the electroweak scale responsible for the W and Z masses and the mass scale of unification. With this requirement it follows that SUSY should be broken at the electroweak scale and therefore SUSY particles are expected to exist at this scale. With this input, one obtains that the gauge coupling constants measured at LEP and other experiments, when evolved via the renormalization group equations to high energies, will join at a scale compatible with proton stability [83].

The simplest SUSY model is the Minimal Supersymmetric Standard Model (MSSM) [84]. This model realizes SUSY in the presence of a discrete R parity (R_p) symmetry. Under this symmetry all standard model particles are even while their partners are odd. As a result of this selection rule SUSY particles are only produced in pairs, with the lightest of them being stable. In the MSSM the Lightest SUSY Particle, LSP for short, is typically a neutralino, for most choices of SUSY parameters. It has been suggested as a candidate for the cold dark matter of the universe and several methods of detection at underground installations have been suggested [85]. However, one should not forget that R parity is postulated ad hoc, without a deep theoretical basis. Moreover there are other ways to explain the cold dark matter via the axion. Last, but not least, hot dark matter is needed in any case, not to mention other existing puzzles in neutrino physics, such as the solar neutrino deficit. From this point of view the emphasis of the simplest MSSM picture would seem exaggerated.

5.2. Broken R–Parity

Unfortunately there is no firm theoretical basis as to how SUSY is realized (if at all) in Nature. Nobody knows the origin of the R parity symmetry. As a matter of fact it could well be broken via tri-linear and bi-linear superpotential couplings. The latter has been shown to be compatible with minimal supergravity with universal boundary conditions at unification [26, 28] as well as the smallness of neutrino masses. This happens for relatively large values of the relevant model effective R–parity violation parameter $\epsilon$ [28]. It provides the simplest reference model for the breaking of R–parity, in the same way as the MSSM provides the simplest phenomenological model for SUSY.

A more satisfactory picture to R–parity violation would be one in which it is conserved at the Lagrangian level but breaks spontaneously through a sneutrino VEV. Keeping the minimal $SU(2) \otimes U(1)$ gauge structure this also implies the spontaneous breaking of lepton number, which is a continuous ungauged symmetry, and therefore the existence of an associated Goldstone boson (majoron). The breaking of R-parity should be driven by isosinglet

\[^{2}\text{For three generations there are three } \epsilon_i, \text{ but here for simplicity we focus only on one.}\]
right-handed sneutrino vacuum expectation values (VEVS) so as to avoid conflicts with LEP observations of the invisible Z width (in this case the majoron is mostly singlet, and does not couple appreciably to the Z). The theoretical viability of this scenario has been demonstrated both with tree-level breaking of the electroweak symmetry and R–parity, as well as in the most attractive radiative breaking approach.

Typically in these models neutrinos have mass. In the conceptually simplest models the origin of neutrino mass is the breaking of R–parity, as in In this case the magnitude of R–parity violating effects is directly correlated with the neutrino mass. However, other consistent possibilities exist.

The model with bi-linear breaking is specially attractive because of its simplicity and because it is the effective truncated version of the more complete models with spontaneous breaking of R–parity.

In the following few sections I will illustrate with some examples the potential of the present and future colliders in testing supersymmetry with spontaneous or bilinear breaking of R–parity under the assumption that neutrinos acquire mass only due to R–parity violation. For simplicity we will refer to these models generically as RPSUSY models. The characteristic feature of these models is that the pattern of R–parity breaking interactions is determined in terms of relatively few new parameters in addition to those of the MSSM (one in the simplest reference model). This allows for a systematic discussion of the potential of new colliders in searching for broken R–parity SUSY signals.

5.2.1. R–Parity Violation at LEP

In the MSSM the usual neutralino pair-production process,

$$e^+ e^- \to Z \to \chi \chi$$  \hspace{1cm} (9)

where \(\chi\) denotes the lightest neutralino, leads to no experimentally detectable signature (other than the contribution to the Z invisible width), as \(\chi\) escapes the apparatus without leaving any tracks. The simplest process that leads to a zen-event topology, with particles in one hemisphere and nothing on the opposite, requires the production of \(\chi\) associated to \(\chi'\), \(Z \to e^+ e^- \to \chi \chi'\), \(\chi'\) being the next-to-lightest neutralino.

In broken R–parity models the \(\chi\) may decay into charged particles, so that eq. (9) can lead to zen-events in which one neutralino decays visibly (leptons and jets) and the other invisibly. The topology is the same as in the MSSM but the corresponding zen-event rates can be larger than in the MSSM and may occur below the threshold for \(\chi'\) production. The missing momentum in these models is carried by the \(\nu_\tau\) or by majorons. Another possibility for zen events in RPSUSY is the decay \(Z \to \chi \nu_\tau\). Since the latter violates R–parity, the rates are somewhat smaller, see. Table 1.

For the sake of illustration we exhibit in Fig. (13) typical values of the branching ratios of neutralinos and charginos, as a function of \(\epsilon\) for \(\mu = 150\) GeV, \(M_2 = 100\) GeV, and \(\tan \beta = 35\). For neutralinos we exhibit its total visible and invisible branching ratios, where we included in the invisible width the contributions coming from the neutrino plus majoron
channel ($\chi \rightarrow \nu J$), as well as from the neutral current channel when the $Z$ decays into a pair of neutrinos ($\chi \rightarrow 3\nu$).

In Fig. (16) we illustrate the sensitivity of LEP experiments to leptonic signals associated to neutralino pair-production at the $Z$ peak in our RPSUSY models. The signal topology used was missing transverse momentum plus acoplanar muon events ($p_T + \mu^+\mu^-$) arising from $\chi\chi$ production followed by $\chi$ decays. The solid line (a) in Fig. (16) is the region of sensitivity of LEP I data of ref. [89] corresponding to an integrated luminosity of $82 \; pb^{-1}$, while (b) corresponds to the improvement expected from including the $e^+e^-\nu$ channel, as well as the combined statistics of the four LEP experiments. The dashed line corresponds to the bilinear model of explicit R–parity violation, allowing $m_{\nu_\tau}$ values as large as the present limit, the dotted one does implement the restriction on $m_{\nu_\tau}$ suggested by nucleosynthesis, and the dash-dotted one is calculated in the model with spontaneous breaking of R–parity (majoron model). The inclusion of semi-leptonic decays and of the updated integrated luminosity already achieved at LEP would substantially improve the statistics and thus the sensitivity to RPSUSY parameters.

The usual chargino pair-production process,

$$e^+e^- \rightarrow Z \rightarrow \chi^+\chi^- \quad (10)$$

may also provide novel signatures which would not be possible in the MSSM, as the neutralinos produced from chargino decays may themselves decay into jets or leptons leading to exotic channels.

Moreover, in $SU(2) \otimes U(1)$ models with spontaneous violation of R–parity the presence of the majoron implies the existence of two–body chargino decays [90],

$$\chi^\pm \rightarrow \tau^\pm + J \quad (11)$$
Figure 16. Limits on $BR(Z \rightarrow \chi\chi)BR(\chi \rightarrow \mu^+ \mu^- \nu)$ compared with the maximum theoretical values expected in different RPSUSY models.

In ref. [91], chargino pair production at LEP II has been studied in supersymmetric models with spontaneously broken $R$–parity. Through detailed signal and background analyses, it was shown that a large region of the parameter space of these models can be probed through chargino searches at LEP II.

The limits on the chargino mass depend on the magnitude of the effective $R$–parity violation parameter $\epsilon$. As $\epsilon \rightarrow 0$ we recover the usual MSSM chargino mass limits, however, for $\epsilon$ sufficiently large, the bounds on the chargino mass can be about 15 GeV weaker than in the MSSM due to the dominance of the two-body chargino decay mode eq. (11). This happens because of the irreducible background from W-pair production with each $W \rightarrow \tau \nu$.

Although the $\nu_\tau$ can be quite relatively heavy in these models, it is consistent with the cosmology critical density [46] as well as primordial nucleosynthesis [58], due to the existence of the majoron which opens new $\nu_\tau$ decay and annihilation channels [48, 50]. The small mass difference between $\nu_e$ and $\nu_\mu$ may lead to an explanation of solar neutrino deficit by resonant $\nu_e$ to $\nu_\mu$ conversions [7]. In this model one may regard the the R–parity violating processes as a tool to probe the physics underlying the solar neutrino conversions [92]. For example, the rates for some RPSUSY rare decays (see section 5.2.3 below) can be used in order to discriminate between large and small mixing angle MSW solutions to the solar neutrino problem [7].

5.2.2. $R$–Parity Violation at LHC It is also possible to find manifestations of $R$–parity violation at the super-high energies available at hadron super-colliders such as the Tevatron and the LHC. If SUSY particles, gluinos and squarks, are pair produced at hadron collisions, their subsequent cascade decays will not terminate at the lightest neutralino but it will further decay. To the extent that this decay is into charged leptons it will give rise to a quite rich pattern of high multiplicity lepton events. Such pattern of gluino cascade decays in RPSUSY...
models was studied in detail in ref. [93]. The conclusion is that multi-lepton and same-sign dilepton signal rates which can be substantially higher than those predicted in the MSSM. This is illustrated in Fig. (18), which shows the branching ratios for various multi-lepton signals (summed over electrons and muons) with the 3-, 4-, 5- and 6-leptons, for \( \tan \beta = 2 \), with other parameters chosen in a suitable way (see ref. [93] for details). We show a) the 3-lepton, b) the 4-lepton, c) the 5-lepton and d) the 6-lepton signal for the MSSM (full line), the majoron-model (dashed line) and the \( \epsilon \)-model (dashed-dotted line). The shaded area will be covered by LEP2. Note, for example, that for \( \mu < 0 \) the 5-lepton signal is much larger in the majoron-model than in the MSSM, giving about 30 to 1200 events per year for an LHC luminosity of \( 10^5 \) pb\(^{-1} \). The 6-lepton signal has a rate up to \( 5 \times 10^{-5} \) in the range \(-300 \text{ GeV} < \mu < -80 \text{ GeV} \) giving 125 events per year. The multi-lepton rates would be even higher in the \( \epsilon \) model.

Figure 17. 95% CL excluded region in RPSUSY models in various analyses (dark areas), as well as the combined excluded region for \( \tan \beta = 2 \), \( \epsilon = 1 \text{ GeV} \), \( \sqrt{s} = 172 \text{ GeV} \), and an integrated luminosity of 300 pb\(^{-1} \).
Although with smaller rates, one also expects in RPSUSY models the single production of the SUSY states in hadron collisions, as has been discussed. In ref. [94] the single production of weakly interacting SUSY fermions (charginos and neutralinos) via the Drell-Yan mechanism was studied.

5.2.3. Rare Decays  If R–parity is broken spontaneously it shows up in the couplings of the W and the Z. As a result there may be rare Z decays with single production of the charginos [90],

$$Z \rightarrow \chi^{\pm} \tau^{\mp}$$  \hspace{1cm} (12)

As mentioned in the RPSUSY models, the magnitude of R parity violation is correlated with the nonzero value of the $\nu_\tau$ mass and is restricted by a variety of experiments. Nevertheless the R parity violating Z decay branching ratios, as an example, can easily exceed $10^{-5}$, well within present LEP sensitivities.

Similarly, the lightest neutralino (LSP) could also be singly-produced as [90]

$$Z \rightarrow \chi \nu_\tau$$  \hspace{1cm} (13)
In models with spontaneous lepton number violation, like majoron models of neutrino mass \[93\], and the SUSY model with spontaneous violation of R–parity \[96\], there are Z decay processes with single photon emission

\[
Z \rightarrow \gamma + H
\]

\[
Z \rightarrow \gamma + J
\]

\[
Z \rightarrow \gamma + J + J
\]

where H is a CP-even Higgs boson and J denotes the associated CP-odd majoron. Since lepton number violation occurs in these models at the weak scale, these processes may have relatively high rates, as seen in Table 1. Their existence would give rise to some rare Z decay signatures which could potentially be observable at the Z peak. In the RPSUSY majoron model \[96\] the first two processes in eq. (14) violate R–parity and are therefore strictly correlated with the $\nu_\tau$ mass. On the other hand in some majoron masses with radiatively induced neutrino mass, like in Fig. \[1\], the branching ratios can be relatively large (possibly accessible at LEP), even though the loop-induced neutrino masses are very small, as required in order to explain the deficit of solar neutrinos.

Another possible signal of the RPSUSY models based on the simplest \(SU(2) \otimes U(1)\) gauge group is rare decays of muons and taus such as \[97\]

\[
\tau \rightarrow \mu + J
\]

\[
\tau \rightarrow e + J
\]

\[
\mu \rightarrow e + J
\]

Such decays would be "seen" as bumps in the final lepton energy spectrum, at half of the parent lepton mass in its rest frame. Again, since in this model the lepton number is broken close to the weak scale it can lead to relatively large rates for single majoron emitting $\mu$ and $\tau$ decays \[97\] compatible with present sensitivities and quite interesting for future tau-charm and B factories \[23\].

Table 5.2.3 summarizes the expectations for rare decay branching ratios in the class of models of interest. It is interesting to note that in the broken R–parity majoron since the processes in eq. \[13\] violate R–parity, their branching ratios are correlated with the $\nu_\tau$ mass. However it was shown in ref. \[97\] that they can be large for moderately small $\nu_\tau$ masses. This illustrates again how the search for rare decays can be a more sensitive probe of neutrino properties than the more direct searches for neutrino masses, and therefore complementary.

5.3. Scalar Sector

Although quite indirect, another possible manifestation of the properties of neutrinos is in the electroweak breaking sector. Many extensions of the lepton sector seek to give masses to
neutrinos through the spontaneous violation of an ungauged U(1) lepton number symmetry, thus implying the existence of a physical Goldstone boson, called majoron [4]. In order to be consistent with the measurements of the invisible Z decay width at LEP the majoron should be (mostly) a singlet under the SU(2) \( \otimes U(1) \) gauge symmetry.

Although the original majoron proposal was made in the framework of the minimal seesaw model, and required the introduction of a relatively high energy scale associated to the mass of the right-handed neutrinos [4], there are many attractive theoretical alternatives where lepton number is violated spontaneously at the weak scale or lower. In this case although the majoron has very tiny couplings to matter and the gauge bosons, it can have significant couplings to the Higgs bosons. As a result the Higgs boson may decay with a substantial branching ratio into the invisible mode [5]

\[ H \rightarrow J + J \] (16)

where \( J \) denotes the majoron. The presence of this invisible decay channel can affect the corresponding Higgs mass bounds in an important way.

The production and subsequent decay of a Higgs boson which may decay visibly or invisibly involves three independent parameters: its mass \( M_H \), its coupling strength to the Z, normalized by that of the SM, \( \epsilon^2 \), and its invisible decay branching ratio. The LEP searches for various exotic channels can be used in order to determine the regions in parameter space that are already ruled out. The exclusion contour in the plane \( \epsilon^2 \) vs. \( M_H \), was shown in Fig. (19) taken from ref. [98].

Another mode of production of invisibly decaying Higgs bosons is that in which a CP even Higgs boson is produced at LEP in association with a massive CP odd scalar [98].
This production mode is present in all but the simplest majoron model containing just one complex scalar singlet in addition to the SM Higgs doublet. As seen in Fig. (20), the cross section for this is typically higher than for the ZH channel, as long as \( m_A \) is not too large. Present limits on the relevant parameters are given in Fig. (21), taken from ref. [98]. In this plot we have assumed \( \text{BR} (H \rightarrow J J) = 100\% \) and a visibly decaying A boson. Similar analysis were made for the case of a high energy linear \( e^+e^- \) collider (NLC) [99], as well as for the LHC [100].

Before concluding let me mention that there are many novel properties of SUSY Higgs bosons, squarks and sleptons that have recently been explored and which may have important implications for the corresponding searches at colliders. For a set of recent references see [101].

5.4. Outlook
1. There are good hints from experiment that neutrinos may be massive, a possibility which is very attractive from the theoretical point of view. This opens the way to many signatures such as neutrino oscillations, neutrino-less double beta decays, possible distortions in beta decay spectra and lepton flavour violating processes.

2. The theoretical attractiveness of supersymmetry justifies the effort devoted to the associated physics and its possible manifestations at present and future particle colliders. So far the negative searches for supersymmetric particles explore little of the relevant region and rely strongly on the *ad hoc* assumption of R–parity conservation.

3. There is a wealth of phenomena that could be associated both to the physics of neutrino mass and to SUSY. They cover a very broad range of energies and experimental situations. In this talk I have considered two examples:

   - models where neutrino masses are generated through lepton number violation at the weak scale
   - R–parity violation as the origin of neutrino mass.

These two classes of schemes are theoretically attractive and lead to a plethora of processes that could be seen in the new generation of particle colliders, from the present LEP II and Tevatron to the future LHC and NLC. R–parity violation provides a hybrid model for neutrino mass generation, combining radiative and seesaw ideas. One finds that $m_{\nu}$ is naturally suppressed and calculable in terms of SUSY parameters and the b-quark Yukawa coupling, while the magnitude of R–parity violation at colliders can be large.

In short, detecting neutrino masses is one of the main challenges in particle physics, with far-reaching implications also for the understanding of fundamental issues in astrophysics and
cosmology. On the other hand, probing Higgs boson physics and searching for SUSY in whatever form is undoubtedly the main goal in the agenda of the next generation of experiments, from elementary particle colliders down to underground experiments and neutrino telescopes. We should be prepared for exciting times where physics beyond the desert will show up!

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