Absolute neutrino masses provide a key to physics beyond the standard model. We discuss the impact of absolute neutrino masses on physics beyond the standard model, the experimental possibilities to determine absolute neutrino masses, and the intriguing connection with the Z-burst model for extreme-energy cosmic rays.

1 Introduction

Solar and atmospheric neutrino oscillations have established solid evidence for non-vanishing neutrino masses and determined mass squared differences $\delta m^2$ and the mixing matrix $U$ with increasing accuracy. The puzzle of the absolute mass scale for neutrinos, though, is still unsolved. In fact it is a true experimental challenge to determine an absolute neutrino mass below 1 eV. Three approaches have the potential to accomplish the task, namely larger versions of the tritium end-point distortion measurements, limits from the evaluation of the large scale structure in the universe, and next-generation neutrinoless double beta decay ($0\nu\beta\beta$) experiments. In addition there is a fourth possibility: the extreme-energy cosmic-ray experiments in the context of the recently emphasized Z-burst model.

This article is organized as follows: in section 2 the specific role of the neutrino among the elementary fermions of the SM is reviewed, as are the two most popular mechanisms for neutrino mass generation. Also, the link of absolute neutrino masses to the physics underlying the standard model is discussed. Section 3 deals with direct determinations of the absolute neutrino mass via tritium beta decay and cosmology. In section 4 $0\nu\beta\beta$ is discussed, which may test very small values of neutrino masses, when information obtained from oscillation studies is input. Section 5 finally deals with the connection of the sub-eV neutrino mass scale and the ZeV energy scale of extreme...
energy cosmic rays in the Z-burst model.

2 Neutrino masses and physics beyond the standard model

The specific role of the neutrino among the elementary fermions of the standard model (SM) is twofold: It is the only neutral particle, and its mass is much smaller than the masses of the charged fermions. Thus it is self-evident that these properties may be related in a deeper theoretical framework underlying the standard model – usually via Majorana mass-generating mechanisms. In the following we comment on the two most popular mechanisms to generate small neutrino masses, namely the see-saw mechanism and radiative neutrino mass generation.

2.1 The see-saw mechanism

The see-saw mechanism is based on the observation, that in order to generate Dirac neutrino masses

$$m_D \nu_L \nu_R$$ (1)

analogous to the mass terms of the charged leptons, the introduction of right-handed neutrinos is required. However a lepton-number violating Majorana mass term for right-handed neutrinos

$$\nu_R M_R (\nu_R)^c$$ (2)

is not prohibited by any gauge symmetry of the standard model. Thus by buying a Dirac neutrino mass term \(m_D\), one inevitably invites mixing with a Majorana mass \(M_R \gg m_D\) which may live a priori at the mass scale of the underlying unified theory. The diagonalization of the general mass matrix yields mass eigenvalues

$$m_\nu \simeq (m_D)^2/M_R \ll m_D$$ (3)

$$M \simeq M_R,$$ (4)

explaining the smallness of the light mass. Though the fundamental scale \(M_R\) is unaccessible for any kind of direct experimental testing, it is obvious from eq. 3 that with information on the low-energy observables \(m_\nu\) and \(m_D\) the “beyond the SM” mass scale of \(M_R\) can be reconstructed. While it turns out to be unrealistic to determine \(m_D\) in the standard model, this option exists indeed in supersymmetry. In the supersymmetric version of the see-saw mechanism, lepton flavor violation (LFV) in the neutrino mass matrix

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\[ \delta \tilde{m}_L^2 \propto Y_\nu Y_\nu^\dagger. \]  

(5)

These LFV soft terms induce large branching ratios for SUSY mediated loop-decays such as \( \mu \to e\gamma \),

\[ \Gamma(\mu \to e\gamma) \propto \alpha^3 \frac{|(\delta \tilde{m}_L)_{12}|^2}{m_S^2} \tan^2 \beta. \]  

(6)

Here \( m_S \) denotes the slepton mass scale in the loop. Thus in the supersymmetric framework it is possible to probe the heavy mass scale \( M_R \) by determining the (light) neutrino mass scale \( m_\nu \) and the (Dirac) Yukawa couplings. This fact is illustrated in fig. 1, where the branching ratio \( \mu \to e\gamma \) is shown as a function of \( M_R \) for a specific mSUGRA scenario, both for a large (lower curve) and small (upper curve) neutrino mass scale.

2.2 Radiative neutrino masses

An alternative mechanism generates neutrino masses via loop graphs at the SUSY scale, in contrast to the tree level generation of charged lepton masses.
Figure 2. Radiative generation of neutrino Majorana masses in $R_p$-violating SUSY.

via the Higgs mechanism (see e.g. 4). In supersymmetry lepton-number violating couplings $\lambda$ and $\lambda'$ may arise if the discrete R-parity symmetry is broken ($R_p$). These couplings may induce neutrino masses via one loop self-energy graphs, see fig. 2. The entries in the neutrino mass matrix, given by

$$m_{\nu_{\alpha\beta}} \simeq \frac{N_c X_{\nu_{\alpha\beta} H}}{16\pi^2} m_d m_d \left[ \frac{f(m_d^2 / m_{\tilde{d}}^2)}{m_{\tilde{d}}^2} + \frac{f(m_{\tilde{d}}^2 / m_d^2)}{m_d^2} \right]$$

are proportional to the products of $R_p$-couplings and depending on the values of superpartner masses. A determination of the absolute neutrinos mass scale would allow one to constrain all entries in the mass matrix, using the smallness of atmospheric and solar $\delta m^2$‘s and unitarity of the neutrino mixing matrix $U$.

In fact, recent bounds on absolute neutrino masses improve previous bounds on $R_p$-couplings by up to 4 orders of magnitude 5. Thus determining the neutrino mass probes physics at heavy mass scales beyond the SM also in the case of radiative generated neutrino masses.

3 Absolute neutrino masses: direct determinations

3.1 Tritium beta decay

In tritium decay, the larger the mass states comprising $\bar{\nu}_e$, the smaller is the $Q$-value of the decay. The manifestation of neutrino mass is a reduction of phase space for the produced electron at the high energy end of its spectrum. An expansion of the decay rate formula about $m_{\nu_e}$ leads to the end point sensitive factor

$$m_{\nu_e}^2 \equiv \sum_j |U_{ej}|^2 m_j^2$$

$$$$
where the sum is over mass states which can kinematically alter the end-point spectrum. If the neutrino masses are nearly degenerate, then unitarity of $U$ leads immediately to a bound on $\sqrt{m_\nu^2} = m_3$. The design of a larger tritium decay experiment (KATRIN) to reduce the present 2.2 eV $m_\nu$ bound is under discussion; direct mass limits as low as 0.4 eV, or even 0.2 eV, may be possible in this type of experiment.

3.2 CMB/LSS cosmological limits

According to Big Bang cosmology, the masses of nonrelativistic neutrinos are related to the neutrino fraction of closure density by $\sum_j m_j = 40 \Omega_\nu h_{65}^2$ eV, where $h_{65}$ is the present Hubble parameter in units of 65 km/s/Mpc. As knowledge of large-scale structure (LSS) formation has evolved, so have the theoretically preferred values for the hot dark matter (HDM) component, $\Omega_\nu$. In fact, the values have declined. In the once popular HDM cosmology, one had $\Omega_\nu \sim 1$ and $m_\nu \sim 10$ eV for each of the mass-degenerate neutrinos. In the cold-hot CHDM cosmology, the cold matter was dominant and one had $\Omega_\nu \sim 0.3$ and $m_\nu \sim 4$ eV for each neutrino mass. In the currently favored $\Lambda$CDM cosmology, there is scant room left for the neutrino component.

The power spectrum of early-Universe density perturbations is processed by gravitational instabilities. However, the free-streaming relativistic neutrinos suppress the growth of fluctuations on scales below the horizon (approximately the Hubble size $c/H(z)$) until they become nonrelativistic at $z \sim m_j/3T_0 \sim 1000 (m_j/eV)$.

A recent limit derived from the 2dF Galaxy Redshift Survey power spectrum constrains the fractional contribution of massive neutrinos to the total mass density to be less than 0.13, translating into a bound on the sum of neutrino mass eigenvalues, $\sum_j m_j < 1.8$ eV (for a total matter mass density $0.1 < \Omega_m < 0.5$ and a scalar spectral index $n = 1$). Previous cosmological bounds come from the data of galaxy cluster abundances, the Lyman $\alpha$ forest, data compilations of the cosmic microwave background (CMB), and peculiar velocities of large scale structure, and give upper bounds on the sum of neutrino masses in the range 3-6 eV. A discussion of possible limits from future supernova neutrino detection is given in [8].

Some caution is warranted in the cosmological approach to neutrino mass, in that the many cosmological parameters may conspire in various combinations to yield nearly identical CMB and LSS data. An assortment of very detailed data may be needed to resolve the possible “cosmic ambiguities”.

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4 Neutrinoless double beta decay

$0\nu\beta\beta$ corresponds to two single beta decays occurring simultaneously in one nucleus and converts a nucleus $(Z,A)$ into a nucleus $(Z+2,A)$. While even the standard model (SM) allowed process emitting two antineutrinos

$$\frac{A}{2}X \rightarrow \frac{A}{2+2}X + 2e^- + 2\bar{\nu}_e$$

(9)

is one of the rarest processes in nature with half lives in the region of $10^{21−24}$ years, more interesting is the search for the neutrinoless mode,

$$\frac{A}{2}X \rightarrow \frac{A}{2+2}X + 2e^-$$

(10)

which violates lepton number by two units and thus implies physics beyond the SM.

The $0\nu\beta\beta$ rate is a sensitive tool for the measurement of the absolute mass-scale for Majorana neutrinos [10]. The observable measured in the amplitude of $0\nu\beta\beta$ is the $ee$ element of the neutrino mass-matrix in the flavor basis (see fig. 3). Expressed in terms of the mass eigenvalues and neutrino mixing-matrix elements, it is

$$m_{ee} = |\sum_i U_{e1}^2 m_i|.$$  

(11)

A reach as low as $m_{ee} \sim 0.01$ eV seems possible with proposed double beta decay projects such as GENIUSI, MAJORANA, EXO, XMASS or MOON. This provides a substantial improvement over the current bound, $m_{ee} < 0.6$ eV.

A recent claim [11] by the Heidelberg-Moscow experiment reports a best fit value of $m_{ee} = 0.36$ eV, but is subject to possible systematic uncertainties.

In the far future, another order of magnitude in reach is available to the 10 ton version of GENIUS, should it be funded and commissioned.

For masses in the interesting range $\gtrsim 0.01$ eV, the two light mass eigenstates are nearly degenerate and so the approximation $m_1 = m_2$ is justified. Due to the restrictive CHOOZ bound, $|U_{e3}|^2 < 0.025$, the contribution of the third mass eigenstate is nearly decoupled from $m_{ee}$ and so $U_{e3}^2 m_3$ may be neglected in the $0\nu\beta\beta$ formula. We label by $\phi_{12}$ the relative phase between $U_{e1}^2 m_1$ and $U_{e2}^2 m_2$. Then, employing the above approximations, we arrive at a very simplified expression for $m_{ee}$:

$$m_{ee}^2 = \left[1 - \sin^2(2\theta_{s\,\odot}) \sin^2 \left(\frac{\phi_{12}}{2}\right)\right] m_1^2.$$  

(12)

The two CP-conserving values of $\phi_{12}$ are 0 and $\pi$. These same two values give maximal constructive and destructive interference of the two dominant terms.
in eq. (11), which leads to upper and lower bounds for the observable $m_{ee}$ in terms of a fixed value of $m_1$:

$$\cos(2\theta_{\text{sun}}) m_1 \leq m_{ee} \leq m_1, \quad \text{for fixed } m_1.$$  \hspace{1cm} (13)

The upper bound becomes an equality, $m_{ee} = m_1$, if $\phi_{12} = 0$. The lower bound depends on Nature’s value of the mixing angle in the LMA solution. A consequence of eq. (13) is that for a given measurement of $m_{ee}$, the corresponding inference of $m_1$ is uncertain over the range $[m_{ee}, m_{ee} \cos(2\theta_{\text{sun}})]$ due to the unknown phase difference $\phi_{12}$, with $\cos(2\theta_{\text{sun}}) \gtrsim 0.1$ weakly bounded even for the LMA solution. This uncertainty disfavors $0\nu\beta\beta$ in comparison to direct measurements if a specific value of $m_1$ has to be determined, while $0\nu\beta\beta$ is more sensitive as long as bounds on $m_1$ are aimed at. Knowing the value of $\theta_{\text{sun}}$ better will improve the estimate of the inherent uncertainty in $m_1$. For the LMA solar solution, the forthcoming Kamland experiment should reduce the error in the mixing angle $\sin^2 2\theta_{\text{sun}}$ to $\pm 0.1$. [4]

5 Extreme energy cosmic rays in the Z-burst model

It was expected that the EECR primaries would be protons from outside the galaxy, produced in Nature’s most extreme environments such as the tori or radio hot spots of active galactic nuclei (AGN). Indeed, cosmic ray data show a spectral flattening just below $10^{19}$ eV which can be interpreted as a new extragalactic component overtakeing the lower energy galactic component; the energy of the break correlates well with the onset of a Larmor radius for protons too large to be contained by the Galactic magnetic field. It was
Figure 4. Schematic diagram showing the production of a Z-burst resulting from the resonant annihilation of a cosmic ray neutrino on a relic (anti)neutrino. If the Z-burst occurs within the GZK zone (∼50 to 100 Mpc) and is directed towards the earth, then photons and nucleons with energy above the GZK cutoff may arrive at earth and initiate super-GZK air-showers.

Further expected that the extragalactic spectrum would reveal an end at the Greisen-Kuzmin-Zatsepin (GZK) cutoff energy of $E_{\text{GZK}} \sim 5 \times 10^{19}$ eV. The origin of the GZK cutoff is the degradation of nucleon energy by the resonant scattering process $N + \gamma \rightarrow \Delta^* \rightarrow N + \pi$ when the nucleon is above the resonant threshold $E_{\Delta^*}$. The concomitant energy-loss factor is $\sim (0.8)^{D/6\text{Mpc}}$ for a nucleon traversing a distance $D$. Since no AGN-like sources are known to exist within 100 Mpc of earth, the energy requirement for a proton arriving at earth with a super-GZK energy is unrealistically high. Nevertheless, to date more than twenty events with energies at and above $10^{20}$ eV have been observed (for recent reviews see [14]).

Several solutions have been proposed for the origin of these EECRs, ranging from unseen Zevatron accelerators ($1$ Zev = $10^{21}$ eV) and decaying supermassive particles and topological defects in the Galactic vicinity, to exotic primaries, exotic new interactions, and even exotic breakdown of conventional physical laws. A rather conservative and economical scenario involves cosmic ray neutrinos scattering resonantly on the cosmic neutrino background (CNB) predicted by Standard Cosmology, to produce Z-bosons [13]. These Z-bosons in turn decay to produce a highly boosted “Z-burst”, containing on average twenty photons and two nucleons above $E_{\text{GZK}}$ (see Fig. 4). The photons and
nucleons from Z-bursts produced within 50 to 100 Mpc of earth can reach earth with enough energy to initiate the air-showers observed at \( \sim 10^{20} \) eV.

The energy of the neutrino annihilating at the peak of the Z-pole is

\[
E_{\nu}^R = \frac{M_Z^2}{2m_j} = 4 \left( \frac{\text{eV}}{m_j} \right) \text{ZeV}.
\] (14)

Even allowing for energy fluctuations about mean values, it is clear that in the Z-burst model the relevant neutrino mass cannot exceed \( \sim 1 \) eV. On the other hand, the neutrino mass cannot be too light or the predicted primary energies will exceed the observed event energies, and the primary neutrino flux will be pushed to unattractively higher energies. In this way, one obtains a rough lower limit on the neutrino mass of \( \sim 0.1 \) eV for the Z-burst model (with allowance made for an order of magnitude energy-loss for those secondaries traversing 50 to 100 Mpc). A detailed comparison of the predicted proton spectrum with the observed EECR spectrum in [14] yields a value of \( m_\nu = 2.34^{+1.29}_{-0.84} \) eV for galactic halo and \( m_\nu = 0.26^{+0.20}_{-0.14} \) eV for extragalactic halo origin of the power-like part of the spectrum.

A necessary condition for the viability of this model is a sufficient flux of neutrinos at \( \gtrsim 10^{21} \) eV. Since this condition seems challenging, any increase of the Z-burst rate that ameliorates slightly the formidable flux requirement is helpful. If the neutrinos are mass degenerate, then a further consequence is that the Z-burst rate at \( E_R \) is three times what it would be without degeneracy. If the neutrino is a Majorana particle, a factor of two more is gained in the Z-burst rate relative to the Dirac neutrino case since the two active helicity states of the relativistic CNB depolarize upon cooling to populate all spin states.

Moreover the viability of the Z-burst model is enhanced if the CNB neutrinos cluster in our matter-rich vicinity of the universe. For smaller scales, the Pauli blocking of identical neutrinos sets a limit on density enhancement. As a crude estimate of Pauli blocking, one may use the zero temperature Fermi gas as a model of the gravitationally bound neutrinos. Requiring that the Fermi momentum of the neutrinos does not exceed mass times virial velocity, one gets the Tremaine-Gunn bound

\[
\frac{\nu_j}{54 \text{ cm}^{-3}} \lesssim 10^3 \left( \frac{m_j}{\text{eV}} \right)^3 \left( \frac{\sigma}{200 \text{ km/s}} \right)^3.
\] (15)

With a virial velocity within our Galactic halo of a couple hundred km/sec it appears that Pauli blocking allows significant clustering on the scale of our Galactic halo only if \( m_j \gtrsim 0.5 \) eV. Free-streaming (not considered here) also works against HDM clustering.
Thus, if the Z-burst model turns out to be the correct explanation of EECRs, then it is probable that neutrinos possess one or more masses in the range $m_\nu \sim (0.1 - 1) \text{ eV}$. Mass-degenerate neutrino models are then likely. Some consequences are:

- A value of $m_{ee} > 0.01 \text{ eV}$, and thus a signal of $0\nu\beta\beta$ in next generation experiments, assuming the neutrinos are Majorana particles.

- Neutrino mass sufficiently large to affect the CMB/LSS power spectrum.

6 Conclusions

The absolute neutrino mass scale is a crucial parameter to learn about the theoretical structures underlying the SM. Information about absolute neutrino masses can be obtained from direct determinations via tritium beta decay or cosmology. More sensitive in giving limits but less valuable for determining the mass scale is neutrinoless double beta decay. The puzzle of EECRs above the GZK cutoff can be solved conservatively with the Z-burst model, connecting the ZeV scale of EECRs to the sub-eV scale of neutrino masses. If the Z-burst model turns out to be correct, neutrino masses in the region of 0.1-1 eV are predicted and degenerate scenarios are favored. In this case positive signals for future tritium beta decay experiments, CMB/LSS studies, and $0\nu\beta\beta$ can be expected.

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