Progress in the understanding of neutrino properties

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Abstract. I briefly summarize neutrino oscillation results and discuss their robustness. I mention recent attempts to understand the pattern of neutrino mixing within various seesaw mechanisms, with or without supersymmetry and/or flavor symmetries. I also mention the possibility of intrinsic supersymmetric neutrino masses in the context of broken R parity models, showing how this leads to clear tests at the LHC.

1. Neutrino Oscillations

Neutrino flavors inter-convert during propagation both in vacuo and in matter. Uncontroversial evidence for this follows from solar and atmospheric, as well as reactor and accelerator neutrino studies [1]. The basic concept needed to describe oscillations is the lepton mixing matrix, the leptonic analogue of the quark mixing matrix. In its simplest $3 \times 3$ unitary form it is given as [2]

$$K = \omega_{23}\omega_{13}\omega_{12}$$

(1)

where each factor is effectively $2 \times 2$ with an angle and a corresponding CP phase. Since current experiments are insensitive to CP violation, we neglect phases, so that oscillations depend only on the three mixing angles $\theta_{12}, \theta_{23}, \theta_{13}$ and on the two squared-mass splittings $\Delta m^2_{21} = m_2^2 - m_1^2$ and $\Delta m^2_{31} = m_3^2 - m_1^2$ associated to solar and atmospheric transitions [3]. To a good approximation, one can also set $\Delta m^2_{21} = 0$ in the analysis of atmospheric and accelerator data, and $\Delta m^2_{31}$ to infinity in the analysis of solar and reactor data. The resulting neutrino oscillation parameters are summarized in Figs. 1 and 2 taken from [4].

![Figure 1](image-url). Neutrino oscillation parameters from a global analysis of the world’s current data [4].
The left and right panels in Fig. 1 give the “atmospheric” and “solar” oscillation parameters, $\theta_{23}$ & $\Delta m_{21}^2$, and $\theta_{12}$ & $\Delta m_{13}^2$, respectively, minimizing with respect to the undisplayed parameters in each case, and including always all relevant data. The dot, star and diamond indicate the best fit points of atmospheric MINOS and global data, as well as solar, KamLAND and global data, respectively. Note the complementarity between data from artificial and natural neutrino sources: reactors and accelerators give the best determination of squared-mass-splittings, while mixings are mainly determined by solar and atmospheric data. Current data slightly prefer a nonzero value of the remaining angle $\theta_{13}$, as seen in Fig. 2. Since the hint is not yet significant, we prefer to interpret this as a weaker bound on $\theta_{13}$:

$$\sin^2 \theta_{13} \leq \begin{cases} 0.060 (0.089) & \text{(solar+KamLAND)} \\ 0.027 (0.058) & \text{(CHOOZ+atm+K2K+MINOS)} \\ 0.035 (0.056) & \text{(global data)} \end{cases}$$

If confirmed by future data, a nonzero $\theta_{13}$ would encourage the search for CP violation in upcoming neutrino oscillation experiments [5, 6]. The expected CP asymmetries are small, as they are suppressed both by $\theta_{13}$ and by the small ratio $\alpha$ of solar/atmospheric squared mass splittings, currently determined as $\alpha = 0.032$, $0.027 \leq \alpha \leq 0.038$ (3$\sigma$).

2. Are oscillations robust?

Many effects may distort the “celestial” neutrino fluxes reaching the detectors, and hence affect the determination of oscillation parameters. In this connection several regular [7–9] and random [10, 11] convective zone solar magnetic field models have been considered. These fields would induce neutrino spin-flavor precession [12–14] and hence modulate the originally produced standard solar model neutrino fluxes. Similarly, the presence of radiative zone random magnetic fields could also induce density fluctuations deep inside the Sun [15] and substantially modify the energy-dependence of the solar neutrino survival probability [16, 17], with a potentially important impact on the determination of oscillation parameters.

Reactor neutrino data from KamLAND have played a crucial role in probing the robustness of solar oscillations against astrophysical uncertainties, such as radiative [18, 19] or convective zone [7–9] magnetic fields. The stringent limits on solar anti-neutrinos from KamLAND lead to strong limits on neutrino magnetic transition moments, especially in the case of turbulent fields [10, 11]. The result is that oscillations constitute the only viable explanation of the data.

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1 Note: the bounds in Eq. (2) are given for 1 dof, while the regions in Fig. 2 (left) are 90% CL for 2 dof.
and of the oscillation solutions allowed by solar data [20], only the large mixing angle solution is consistent with the spectrum measurements at KamLAND [21].

In Standard Model language neutrino mass is described by an effective dimension-five operator, shown in the left panel in Fig. 3.

Most neutrino mass generation schemes also induce sub-weak strength ($\sim \varepsilon G_F$) dimension-6 operators depicted in the right panel in Fig. 3. They can be of two types: flavour-changing (FC) and non-universal (NU). For example, the presence of NSI is expected low-scale seesaw models, such as the inverse [22–24] and the linear [25] seesaw models, where the non-unitary piece [2] of the lepton mixing matrix can be sizeable, hence the induced non-standard interactions. Relatively sizable NSI strengths may also be induced in models with radiatively induced neutrino masses [26–29]. The strength of the NSI operators will play an important role in elucidating the origin of neutrino mass, as it will help discriminate between high and low-scale schemes.

As noted in Ref. [30], current determination of solar neutrino parameters is not yet robust against the existence of non-standard neutrino interactions. In fact, in the presence of such interactions, there is a new “dark side” solution [30], as shown in the left panel in Fig. 4 taken from Ref. [31]. This solution is almost degenerate with the usual one, and survives the inclusion of reactor data. As seen in the right panel in Fig. 4, the lines are compared with the experimental rates for the pp neutrino flux (from the combination of all solar experiments), the 0.862 MeV $^7$Be line (from Borexino) and two estimated values of the $^8$B neutrino flux from Borexino and SNO (third phase). Distinguishing between these solutions poses a challenge for low energy solar neutrino experiments.

As already noted, the slight tension between current solar and KamLAND data is alleviated allowing for a non-zero value of the mixing angle $\theta_{13}$. Likewise, non-standard flavor-changing interactions may also alleviate this tension and hence constitute a source of confusion in the determination of $\theta_{13}$. Indeed, a full three-flavor analysis of solar and KamLAND data in the

Figure 3. Neutrino mass and corresponding NSI operators.

Figure 4. Left: normal and “dark-side” solutions in the presence of non-standard neutrino interactions (NSI). Right: solar neutrino survival probabilities averaged over the $^8$B neutrino production region for three reference points, see Ref. [31] for details. Vertical bars correspond to experimental errors, while horizontal ones indicate the energy range relevant for each experiment.
presence of flavor-changing NSI [32] reveals a degeneracy between $\theta_{13}$ and the vectorial coupling $\epsilon_{13}$ characterizing the non-standard transitions between $\nu_e$ and $\nu_\tau$ in the forward scattering process with d-type quarks.

Non-standard neutrino interactions may also be probed at high energies. One can show that already even a small residual NSI strengths may have dramatic consequences for the sensitivity to $\theta_{13}$ at a neutrino factory [33, 34]. Improving the sensitivities on NSI constitutes a necessary step and opens a window of opportunity for neutrino physics in the precision age. The presence of NSI also leads to novel effects in supernova neutrino propagation [35], for example, the possibility of internal resonant neutrino conversions even in the absence of neutrino masses [36–38].

In contrast to the “solar” sector, thanks to the large statistics of atmospheric data over a wide energy range, the determination of $\Delta m^2_{31}$ and $\sin^2 \theta_{\text{ATM}}$ is hardly affected by the presence of NSI, at least within the 2–neutrino approximation [39]. Future neutrino factories will substantially improve this bound [40].

3. How do neutrinos get mass?

Neutrino oscillations provide the first sign of physics beyond the Standard Model (SM). Pinning-down the ultimate origin of neutrino mass remains a challenge. Apart from any detailed mass generation mechanism, since they carry no electric charge they should have Majorana-type masses [2]. This is indeed what happens in specific neutrino mass generation schemes, markedly different from those of charged fermions. The latter come in two chiral species and get mass linearly in the electroweak symmetry breaking vacuum expectation value (vev) $\langle \Phi \rangle$ of the Higgs scalar doublet. In contrast, as illustrated in the left panel in Fig. 3, neutrino masses arise as an effective lepton number violating dimension-five operator $O \equiv \lambda L \Phi \Phi$ (where $L$ denotes a lepton doublet) [43]. It follows that a simple way to account for the smallness of neutrino masses, is that the coefficient characterizing the strength of $O$ is suppressed either by a high-scale $M_X$ in the denominator or, alternatively, it may involve a low-mass-scale in the numerator. The search of NSI will help discriminate which of the two pathways is chosen by nature.

Gravity is often argued to break global symmetries [44, 45], and would induce Weinberg’s operator $O$, with $M_X$ identified to the Planck scale $M_P$. The resulting masses are too small, implying the need for physics beyond the SM, typically at a high scale below $M_P$ [46]. Alternatively $O$ may be suppressed by small scales, Yukawa couplings and/or loop-factors [47]. This way one may have tree level, radiative, and hybrid mechanisms, all of which may have high- or low-scale realizations. Depending on whether it is gauged or not, spontaneous lepton-number violation implies either an extra neutral gauge boson or a Nambu-Goldstone boson coupled to neutrinos. One may construct models of both types. However the most basic and general seesaw description is in terms of the $SU(3) \otimes SU(2) \otimes U(1)$ gauge structure [2]. In such a framework the relevant scale can either be large or small, depending on model details, with a fair chance that the origin of neutrino mass may be probed at accelerators like the LHC.

3.1. Minimal seesaw

Here the operator $O$ is induced by the exchange of heavy states with masses close to the “unification” scale, as indicated in Fig. 3 for a review see Ref. [47]. In type-I or type-III seesaw the exchanged states are fermions singlets or triplets, while Type-II seesaw is induced by triplet scalar exchange. The “complete seesaw” has both fermion and scalar intermediate states and its phenomenology has been thoroughly studied in [2] including the general seesaw diagonalization method [48]. The latter is obtained using the hierarchy of vevs $v_3 \ll v_2 \ll v_1$ needed in order to account naturally for the small neutrino masses and which follows by minimizing the scalar potential.

Lepton number violation may also show up at higher dimension [41, 42].
3.2. “Non-minimal” seesaw

The seesaw is not a model but a general paradigm. One may implement it with different gauge groups and multiplet contents, with gauged or ungauged B-L, broken explicitly or spontaneously, at a high or at a low energy scale, with or without supersymmetry. An attractive class of extended seesaw schemes employs extra gauge singlet neutrinos in addition to the $\nu^c$ which are present in the 16 of $SO(10)$ group [22]. One may implement such schemes with low-scale breaking of B-L [25,49], leaving a TeV-scale $Z'$ [50]. However, whatever the symmetry is, it must ultimately break to $SU(3) \otimes SU(2) \otimes U(1)$, hence the most general seesaw description must be formulated at this level [2]. Such low-energy description is crucial in accurately describing low-scale variants of the seesaw mechanism, where the basic lepton-number-violating parameter may be naturally small and calculable due to supersymmetric renormalization group evolution [23].

3.3. Radiative schemes

The operator $O$ may arise as loop effects [26,27], with no need for a large scale. In this case its coefficient $\lambda$ is suppressed by small loop-factors, Yukawa couplings and possibly by a small scale parameter characterizing the breaking of lepton number, leading to naturally small neutrino masses. Like low-scale seesaw models discussed above, radiative schemes typically have new TeV states opening the door to phenomenology at the LHC. [51].

3.4. R parity violation

The origin of neutrino masses may be intrinsically supersymmetric in models where the R parity symmetry breaks [52–54]. This may happen spontaneously, driven by a nonzero vev of an $SU(3) \otimes SU(2) \otimes U(1)$ singlet sneutrino [55–57], leading to effective bilinear R-parity violation [58,59]. This provides the minimal way to break R parity in the Minimal Supersymmetric Standard Model [59]. The induced neutrino mass spectrum is hybrid, with one scale (typically the atmospheric) induced by neutralino-exchange at the tree level, and the other scale (solar) induced by calculable one-loop corrections [60].

Unprotected by any symmetry, the lightest supersymmetric particle (LSP) decays, typically inside detectors at the Tevatron or the LHC [60,61]. The LSP decay-length is a measure of the neutrino mass and can be experimentally resolvable, leading to a displaced vertex [62,63]. More strikingly, its decay properties correlate with the neutrino mixing angles. Indeed, the LSP decay pattern is predicted by the low-energy measurement of the atmospheric angle $\theta_{23}$ [64–66]. Such a prediction will be clearly tested at the LHC. Similar correlations hold in schemes based on other supersymmetry breaking mechanisms and, correspondingly, featuring other states as LSP [67].
4. Understanding lepton mixing with flavor symmetries

As seen above current neutrino oscillation data indicate solar and atmospheric mixing angles which are unexpectedly large when compared with quark mixing angles. This places a challenge to the understanding of the flavor problem in unified schemes where quarks and leptons are related. It has been noted that the neutrino mixing angles are approximately tri-bi-maximal [68].

There have been many schemes suggested in the literature in order to reproduce this pattern using various discrete flavor symmetry groups containing mu-tau symmetry, e.g. [69–78]. One expects the flavor symmetry to be valid at high energy scales. Deviations from the tri-bi-maximal ansatz [79] may be calculable by renormalization group evolution [80–82].

A specially simple ansatz is that, as a result of a given flavor symmetry such as A4 [69, 70], neutrino masses unify at high energies $M_X$ [83], the same way as gauge couplings do in the presence of supersymmetry [84]. Such quasi-degenerate neutrino scheme predicts $\theta_{23} = \pi/4$ and $\theta_{13} = 0$, leaving the solar angle $\theta_{12}$ arbitrary and the Dirac phase is maximal [72]. One can show that lepton and slepton mixings are related and that at least one slepton lies below 200 GeV, within reach of the LHC. The lower bound on the absolute Majorana neutrino mass scale $m_0 > \sim 0.3$ eV ensures that the model will be probed by future cosmological tests and $\beta\beta_0\nu$ searches. Expected rates for LFV processes $\text{BR}(\mu \rightarrow e\gamma) > \sim 10^{-15}$ and $\text{BR}(\tau \rightarrow \mu \gamma) > 10^{-9}$ typically lie within reach of upcoming experiments. Note that flavor symmetries, such as our A4, may also be implemented in low-scale seesaw schemes, both type-I [85] and type-III [24], leading to different neutrino mass spectra.

5. Lepton flavor violating (LFV) effects

The unequivocal evidence that neutrinos oscillate suggests that, at some level, flavor violation should also show up as transitions involving the charged leptons, since these are their electroweak doublet partners. There are two basic mechanisms: (i) neutral heavy lepton exchange [86–88] and (ii) supersymmetry exchange [89–91]. Due to the sizeable admixture of right-handed neutrinos in their leptonic charged current, low-scale seesaw schemes induce potentially large LFV rates [86, 87] and CP violating processes as well [93, 94]. This can happen in the absence of supersymmetry and in the limit of massless neutrinos, hence their magnitude is unrestricted by the smallness of neutrino masses. In Fig. 6 we display $\text{Br}(\mu \rightarrow e\gamma)$ versus the small lepton number violating (LNV) parameters $\mu$ and $v_L$ for two different low-scale seesaw models, the inverse and the linear seesaw, respectively. Clearly the LFV rates are sizeable in both cases. Similarly one can see that in low-scale seesaw models the nuclear $\mu^–e^–$ conversion rates lie within planned sensitivities of future experiments such as PRISM [95]. Note that in type-III
versions of such low-scale seesaw schemes [24], the TeV RH neutrinos would not only induce
LFV processes but also be copiously produced at the LHC.

In contrast, barring fine-tunings, high-scale seesaw models require supersymmetry in order
to have sizeable LFV rates. Here lepton flavour violation is expected to show up in the most
direct way in the production of supersymmetric particles at the LHC, as seen in Fig. 7.

![Figure 7](image)

**Figure 7.** LFV rate for $\mu$-$\tau$ lepton pair production from $\chi_2^0$ decays versus $M_{1/2}$ for the indicated
$m_0$ values, assuming minimal supergravity parameters: $\mu > 0$, tan $\beta = 10$ and $A_0 = 0$ GeV, for
type-I (left) and for type-II seesaw (right). Here $\lambda_1 = 0.02$ and $\lambda_2 = 0.5$ are Type-II seesaw
parameters, and we imposed the constraint $\text{Br}(\mu \rightarrow e + \gamma) \leq 1.2 \cdot 10^{-11}$, from Ref. [96].

Both supersymmetric and RH neutrino contributions to lepton flavour violation exist in
supersymmetric seesaw schemes, and their interplay depends on the seesaw scale [97].

6. Lepton number violation and $\beta\beta_0$\nu

The Dirac or Majorana nature of neutrinos is manifest only through the observation of
LNV processes, such as $\beta\beta_0$\nu [2] whose current status and perspectives was reviewed by
Schoenert [98]. A nonzero rate for $\beta\beta_0$\nu implies that, within a gauge theory, at least one neutrino
gets a Majorana mass, this argument is known as the "black-box" theorem [99], illustrated in
Fig. 8 and recently discussed in [100].

![Figure 8](image)

**Figure 8.** Neutrino mass mechanism for $\beta\beta_0$\nu (left), and black box theorem (right) [99].

Given the neutrino oscillation data it follows that light Majorana neutrino exchange will
induce $\beta\beta_0$\nu and the corresponding rate is a measure of the absolute scale of neutrino
mass, complementary to the one probed in beta decay studies [101], and cosmological
observations [102]. It is also sensitive to the Majorana CP violation [103] which drops out of
oscillations. Using current neutrino oscillation parameters and state-of-the-art nuclear matrix
elements [104] one can determine the average mass parameter $\langle m_\nu \rangle$ characterizing the neutrino
exchange contribution to $\beta\beta_0$\nu.

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3 Electromagnetic interactions of neutrinos can also probe their Majorana nature.
Inverted hierarchy implies a generic lower bound for the $\beta \beta$ amplitude, while for normal hierarchy neutrino spectra the three neutrinos can interfere destructively, so that no generic lower bound exists. Specific flavor models may provide a lower bound for $\beta \beta$, even with normal hierarchy, as discussed in [85] [105, 106]. Quasi-degenerate neutrinos [69, 70] give the largest possible $\beta \beta$ signal.

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