Neutrino mass from higher than d=5 effective operators with discrete symmetries

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Abstract. In beyond Standard Model physics scenarios one often has to extend the Higgs sector and introduce additional (discrete) symmetries. In such models it is possible that the Weinberg operator is forbidden and the leading contribution to neutrino mass is generated by operators with mass dimension $d > 5$. These higher-dimensional operators are stronger suppressed and therefore the new physics scale can be as low as the TeV scale. With this approach it is possible to test neutrino mass models at colliders, such as the LHC. We will demonstrate how these operators can be systematically studied and present an exemplary implementation of such a model and its phenomenological consequences. We also show how it can be embedded into an SU(5) GUT framework. Finally we present a specific realization of this model where we can have tri-bimaximal mixing with large deviations due to a softly broken symmetry.

1. Introduction
With the discovery of a Higgs-like particle [1, 2] the last missing part of the Standard Model (SM) seems to be finally found. After the current upgrade, the Large Hadron Collider (LHC) will be able to further test new physics models. Which one of these models may be realized in nature is so far unclear. They share, however, often some common features. In many scenarios we find, for example, an extended Higgs sector, such as in supersymmetric models (SUSY) or generic Two Higgs Doublet Models (THDMs). Often also new fundamental discrete symmetries are introduced, particularly in the context of flavor models. The recent evidence of a non-zero neutrino mixing angle $\theta_{13}$ [3, 4] is especially interesting in this regard.

In literature models that connect neutrino physics to the TeV scale have been widely discussed lately. Examples are radiative mass generation [5–10] or models with a small lepton number violating contribution such as the inverse seesaw mechanism [11] or SUSY with R-parity violation [12–21]. It is possible in models with additional scalars and discrete symmetries that neutrino masses are generated by higher than $d = 5$ effective operators [6, 22–33]. In the following, we will study these operators in the framework of SUSY [34] and an SU(5) grand unified theory (GUT) [35] and will show that this approach is another possibility for neutrino masses from physics at the TeV scale.

2. Neutrino masses and effective field theory
The usual model to describe the smallness of the neutrino masses is the seesaw mechanism, where – in case of the type-I seesaw – heavy right handed singlets are added. These right
handed neutrinos have Yukawa-like interactions with the left-handed SM neutrinos and masses close to the GUT scale. Since this scale is far beyond the scale of neutrino physics, one can integrate out the right handed neutrinos in order to obtain an effective interaction, the Weinberg operator [36],

\[ \mathcal{L}^{d=5}_{\text{eff}} \propto \frac{1}{\Lambda_{\text{NP}}} (\bar{L} \tau^2 H) (H^\dagger \tau^2 L). \]

After electroweak symmetry breaking (EWSB) this operator will generate an effective Majorana neutrino mass term \( m_\nu \propto v^2 / \Lambda_{\text{NP}} \), where \( v \) is the vacuum expectation value (VEV) of the Higgs field and \( \Lambda_{\text{NP}} \) is the new physics scale. In case of the type-I seesaw, \( \Lambda_{\text{NP}} \) is the mass of the right-handed neutrinos. In general, one can describe the low-energy effects of any fundamental theory by a tower of effective operators, which is added to the SM Lagrangian as

\[ \mathcal{L} = \mathcal{L}_{\text{SM}} + \mathcal{L}^{d=5}_{\text{eff}} + \mathcal{L}^{d=6}_{\text{eff}} + \cdots \quad \text{with} \quad \mathcal{L}^{d}_{\text{eff}} \propto \frac{1}{\Lambda_{\text{NP}}} \mathcal{O}^d. \]

If the Weinberg operator is forbidden by a symmetry, the leading contribution to neutrino mass will be generated by an operator with dimension \( d > 5 \). As a consequence, the effective neutrino mass will be suppressed by higher powers of the new physics scale \( \Lambda_{\text{NP}} \), which therefore can be at much lower energies. This opens the window to test the mechanism generating neutrino masses at collider experiments such as the LHC.

As pointed out by the authors of Ref. [28], it is not possible in the SM to have a leading contribution to neutrino mass from a generic higher-dimensional operator, if the Weinberg operator is forbidden. This is due to the fact that all possible operators contributing to neutrino mass are of the type \((LLHH)(H^\dagger H)^n\) and \((H^\dagger H)^n\) is a singlet under any symmetry. Therefore, we want to focus on models which have an additional Higgs doublet such as THDMs or the Minimal Supersymmetric Standard Model (MSSM). We also consider models with an additional scalar singlet such as the Next-to Minimal Supersymmetric Standard Model (NMSSM). Due to the requirement of SUSY invariance, the number of possible operators in supersymmetric models is limited compared to the THDM. The general form of operators leading to neutrino mass in SUSY is

\[ \mathcal{L}^{d=5+2n+m}_{\text{eff}} = \frac{1}{\Lambda^{d-4}} (LLHH_uH_u)(H_uH_d)^nS^m, \quad n, m = 1, 2, 3, \ldots \]

In table 1 we show a list of all operators up to \( d = 9 \) that contribute to neutrino mass in the framework of the NMSSM. The introduction of a discrete symmetry has several consequences: First, the usual terms from the NMSSM superpotential must be invariant under this symmetry. Furthermore we want the \( d = 5 \) operator to be forbidden by this symmetry. Respecting these constraints, we have listed the charge of the higher-dimensional operators in table 1 with regard to the NMSSM fields. As one can see from the table, it is not possible to have certain operators as leading contribution to neutrino mass, since they always will have the same charge as a lower-dimensional operator. As indicated in table 1, one can show that this applies to all operators with \( d > 7 \) that include the scalar singlet \( S \) [34].

3. A d=7 Example

There are three possible tree-level models, which after integrating out the heavy particles lead to the Weinberg operator: the type-I, type-II and type-III seesaw mechanism [5]. In the same way, it is possible to decompose the higher-dimensional operators. The operators #1, #2 and #4 of table 1, for example, have been studied in Ref. [24].
Table 1. Effective operators generating neutrino mass in the NMSSM up to $d = 9$. Taken from Ref. [34].

| Op. # | Effective interaction | Charge | Same as |
|-------|-----------------------|--------|---------|
| $d = 5$ | $LLH_uH_u$ | $2q_L + 2q_{H_u}$ | |
| $d = 6$ | $LLH_uH_uS$ | $2q_L + q_{H_u} - q_{H_d}$ | |
| $d = 7$ | $LLH_uH_uH_u^2H_u$ | $2q_L + 3q_{H_u} + q_{H_d}$ | |
| $d = 8$ | $LLH_uH_uSS$ | $2q_L - 2q_{H_d}$ | |
| $d = 9$ | $LLH_uH_uH_u^2H_uH_u^2$ | $2q_L + 4q_{H_u} + 2q_{H_d}$ | \#1 |
| $d = 9$ | $LLH_uH_uH_u^2H_u^2SS$ | $2q_L + q_{H_u} - q_{H_d}$ | \#2 |
| $d = 9$ | $LLH_uH_uSSSS$ | $2q_L + q_{H_u} - q_{H_d}$ | \#2 |

As has been discussed in more detail in Ref. [34] operator \#3 can have some interesting phenomenological consequences at the LHC. In order to forbid the Weinberg operator, we choose the following charges for the (N)MSSM fields under a $Z_3$ symmetry

$$q_{H_u} = 0, \quad q_{H_d} = 1, \quad q_L = 1, \quad (qs = 2).$$

One particularly interesting decomposition of this $d = 7$ operator is given by the superpotential

$$W = W_{\text{quarks}} + Y_u e^\xi \hat{L} \cdot \hat{H}_d - Y_N \hat{N} \hat{L} \cdot \hat{H}_u + \kappa_1 \hat{N}' \hat{N}' \cdot \hat{H}_u + \kappa_2 \hat{N}' \hat{N}' \cdot \hat{H}_u + m_N \hat{N} \hat{N}'$$

$$+ m_0 \xi^i \cdot \hat{N} + \mu \hat{H}_u \cdot \hat{H}_d,$$

where $\xi$ and $\xi'$ are SU(2) doublets, whereas $N$ and $N'$ are singlets under the SM gauge group. From this superpotential one derives the following mass matrix for the neutral fermions

$$M_f^0 = \begin{pmatrix} 0 & Y_N v_u & 0 & 0 \\ Y_N^T v_u & 0 & m_N^T & 0 \\ 0 & m_N & 0 & \kappa_1 v_d \\ 0 & \kappa_1^T v_d & 0 & -m_\xi \\ 0 & \kappa_2^T v_u & -m_\xi & 0 \end{pmatrix}$$

in the basis $f^0 = (\nu, N, N', \xi^0, \xi'^0)$, where $v_u, v_d$ are the VEVs of the neutral Higgs components. By diagonalizing this matrix one obtains the corresponding mass eigenstates $n_i$.

The effective neutrino mass is after integrating out the heavy fields

$$m_\nu = v_u^3 v_d^2 Y_N^2 \frac{\kappa_1 \kappa_2 m_N^2}{m_\xi m_N^2}.$$ 

Assuming Yukawa couplings of $O(10^{-3})$, the masses of the heavy particles must indeed be at the TeV range in order to reproduce the observed masses for the light neutrinos.

4. Phenomenological implications at the LHC

Due to the small Yukawa couplings of the singlets $N$ and $N'$, their production cross section will be small. The SU(2) doublets $\xi$ and $\xi'$ instead, which couple to the weak gauge bosons, can be
produced at the LHC in Drell-Yan processes [34]. For $m_N = 200 \text{GeV}$, one obtains cross sections $\sigma(pp \to \xi^\pm \xi^0)$ of about 122 fb (417 fb) for 7 (14) TeV. The doublets will decay dominantly into a neutrino mass eigenstate $n_i$ plus a weak gauge boson or one of the Higgs components. Due to the small mixing between the light and heavy neutrino mass eigenstates the doublets will have small decay width. The neutral mass eigenstates $n_i$ can decay into $W^{\pm}l_i^\pm$ and $H^\pm l_i^\pm$ or into neutrinos plus $Z$ or neutral Higgs components with decay width of $O(10^{-5} \ldots 1) \text{keV}$. Therefore these particles can have sizable decay lengths between $100 \mu\text{m}$ and several millimeters, which helps to identify the particles of our model and to suppress the backgrounds. We obtain results as large as several fb for the lepton number conserving processes $pp \to l^+ l^- W^\pm$, where lepton flavor can be violated. For high enough luminosities it is possible to observe them. The following processes, which violate lepton number by two units, can be used to test the Majorana nature of the $n_i$: \[ pp \to \ell^+ \ell^- W^- \] \[ pp \to \ell^+ \ell^- W^- W^-, \ell^- \ell^- W^+ W^+ \] For the process in Eq. (8) we have strong suppression, since the neutral mass eigenstates $n_4$ and $n_5$ form a pseudo-Dirac pair. As a consequence, the cross section, which is proportional to $m_{n_5}^2 - m_{n_4}^2 \approx O(m_N^2)$, will be tiny. For the lepton number violating processes in Eq. (9) we find on the other hand that the cross sections are larger than one might expect. This is due to the fact that $\xi$ and $\xi'$ can be regarded as vector-like representation of SU(2), which leads to some non-vanishing contributions to the LNV cross section. These are, however, at the limit of observability.

5. GUT completion

As discussed in Ref. [35], the introduction of additional particles that are charged under the SM gauge group modifies the running of the gauge couplings. The model presented in the previous section will spoil unification unless we add complete SU(5) multiplets to the (N)MSSM. Therefore we have to embed all charged particles in complete SU(5) multiplets: \[ \tilde{5}_M = \left( \begin{array}{c} (d_R)^c \\ L \end{array} \right), \quad \tilde{5}_{\xi'} = \left( \begin{array}{c} d^{c'} \\ \xi' \end{array} \right), \quad 5_{\xi} = \left( \begin{array}{c} d'' \\ \xi \end{array} \right). \] (10)

Also the Higgs fields obtain additional colored components, which can be integrated out. As pointed out in Ref. [35], this choice of mediators is the only decomposition of the $d = 7$ operator that can be the leading contribution to neutrino mass and guarantees perturbativity in an SU(5) GUT.

We obtain the most interesting results in a model based on the NMSSM. The most general SU(5) invariant realization of this model is given by the superpotential
\[ W = y_1 N \tilde{5}_5 H_\tau + y_2 N \tilde{5}_{\xi'} H_5 + y_3 N \tilde{5}_M H_5 + y'_1 N' \tilde{5}_5 H_\tau + y'_2 N' \tilde{5}_{\xi'} H_5 + y'_3 N' \tilde{5}_M H_5 + \lambda_{\xi} S \tilde{5}_M 5_{\xi} + \lambda_{\xi'} S \tilde{5}_{\xi'} 5_{\xi} + \lambda_{\xi''} S \tilde{5}_{\xi''} 5_{\xi} + \lambda_{N} S N' N + \lambda_{N' N} S N' N' + \lambda_{\xi N} S N N + \lambda_{\xi' N} S N' N' + \lambda_{\xi'' N} S N'' N' + y_d \tilde{5}_M 10 H_\tau + y'_d \tilde{5}_{\xi'} 10 H_\tau + y_u 10 10 H_5. \] (11)

In this scenario the mediator mass scale is determined by $\langle S \rangle$, which is of order 1 TeV and results in the required mass scale for the light neutrinos. Furthermore the NMSSM avoids some problems of the $\mu$-term of the NMSSM (see also Refs. [28, 34, 35]). As discussed in Ref. [35], however, it is not possible to choose a discrete symmetry to forbid the Weinberg operator, since we obtain now effective operators of the type \[ \frac{1}{\langle S \rangle} LLH_u H_u \quad \text{and} \quad \frac{1}{\langle S \rangle^3} (LLH_u H_d)(H_u H_d) \],
The (S) breaks any discrete symmetry under which it is charged.

We therefore consider a model with an additional $S'$ and the superpotential

$$W = y_\tilde{\nu} 5 M H_5 + y'_1 5'_5 H_5 + y'_2 N 5 H_5 + y'_3 N' 5 H_5 + \lambda_S S' H_5 H_5 + \kappa S^3 + \kappa' S'^3 +$$

where $y_\tilde{\nu}$ breaks any discrete symmetry under which it is charged.

We therefore consider a model with an additional $S'$ and the superpotential

$$W = y_\tilde{\nu} 5 M H_5 + y'_1 5'_5 H_5 + y'_2 N 5 H_5 + y'_3 N' 5 H_5 + \lambda_S S' H_5 H_5 + \kappa S^3 + \kappa' S'^3 +$$

We now charge the fields as listed in table 2. Applying these charges to Eq. (12) one sees that the interactions in the first two lines respect the discrete symmetry, and the ones of the last line break it softly. These terms not only allow for a sizable VEV of $S'$ but also generate a $d = 5$ contribution to neutrino mass after integrating out $\xi$ and $\xi'$:

$$m_\nu^{d=5} = \frac{y_3 y'_3 y''_3}{(S')}.$$  

If all soft breaking couplings are of order $10^{-8}$ or smaller this contribution to neutrino mass will be sub-leading. If the soft coupling $y_3$, however, is close to this upper limit, one can, besides the leading $d = 7$ neutrino mass operator, have sizable corrections from the $d = 5$ operator of Eq. (13). Suppose that the $d = 7$ operator has a tri-bimaximal flavor structure, where (for the sake of simplicity) the charged lepton mass matrix is assumed to be diagonal. Then the deviation from tri-bimaximal mixing can originate in the soft breaking terms in Eq. (12).

In the following, we will show an example from Ref. [35], where the $d = 7$ contribution has tri-bimaximal structure but corrections from the soft-breaking $d = 5$ operator lead to non-zero $\theta_{13}$. We assume a flavor structure where we have three generations of the singlet fields $N$ and $N'$, and the following parametrization of the coupling constants

$$y'_1 = \tilde{y}_1 \begin{pmatrix} 0 \\ \rho \end{pmatrix}, \quad y'_2 = \tilde{y}_2 \begin{pmatrix} 0 \\ -1 \rho \end{pmatrix}, \quad y'_3 = \tilde{y}_3 \begin{pmatrix} \sqrt{2/3} & 1/\sqrt{3} & 0 \\ -1/\sqrt{6} & 1/\sqrt{3} & -1/\sqrt{2} \\ -1/\sqrt{6} & 1/\sqrt{3} & 1/\sqrt{2} \end{pmatrix},$$

where $\tilde{y}_1, \tilde{y}_2$ and $\tilde{y}_3$ are numerical (scalar) parameters, $\rho = \sqrt{m_3/m_2}$ and the mass matrices for $N$ and $N'$ are diagonal. This choice of parameters will generate a tri-bimaximal mass matrix for the neutrinos at $d = 7$.

We obtain now a possible structure of the soft-breaking coupling $y'_3$ (using the parametrization from Ref. [37] for deviations from tri-bimaximal mixing) as

$$y'_3 \simeq \tilde{y}_3 r \begin{pmatrix} 0 & 0 & -\frac{1}{\sqrt{2}} \\ \frac{1}{\sqrt{6}} & \frac{1}{2\sqrt{3}} & 0 \\ -\frac{1}{\sqrt{6}} & -\frac{1}{2\sqrt{3}} & 0 \end{pmatrix} \rho^2 - \begin{pmatrix} 0 & 0 & \frac{3\sqrt{2}}{2} \\ \frac{1}{2\sqrt{3}} & \frac{3\sqrt{2}}{2} & 0 \\ 0 & 0 & \frac{3\sqrt{2}}{2} \end{pmatrix},$$

with $O(\tilde{y}_3) = 10^{-8}$. This results in a correction to the tri-bimaximal mass matrix from the $d = 7$ operator that leads to $\theta_{13} > 0$.  

Table 2. Charges for the fields of the model defined in Eq. (12). Taken from Ref. [35].

| Multiplet  | $5_M$ | $H_5$ | $H'_5$ | $N$ | $N'$ | $5_\xi$ | $5_\xi'$ | $10$ | $S$ | $S'$ |
|-----------|-------|-------|-------|-----|------|--------|---------|-----|-----|------|
| $Z_3$ charge | 1     | 1     | 1     | 1   | 2    | 0      | 0       | 1   | 1   | 0    |
Another aspect of the GUT completion of this model is that the promotion of the SM SU(2) doublets to 5-plets of SU(5) requires the introduction of additional d-quarks as indicated in Eq. (10). Renormalization group effects imply that the masses of these quarks are as large as several TeV, which makes them potentially interesting for LHC studies. Furthermore, in our model these particles are stable due to the same symmetry that suppresses the Weinberg operator. Heavy stable quarks are forbidden due to bounds from Big Bang nucleosynthesis and searches for heavy nuclei. In the case of a softly broken symmetry, however, the heavy quarks can decay via the two-body decays

\[ D' \rightarrow H^- u, \ H^0 d \ . \]  

(16)

with life-times for the \( D' \) as small as \( 10^{-10} \text{ s} \) to \( 10^{-13} \text{ s} \) [35].

6. Conclusions
We have demonstrated that in new physics scenarios with an extended Higgs sector and additional symmetries neutrino masses can be generated by operators with mass dimension \( d > 5 \). The additional suppression of these operators connects neutrino physics to the TeV scale. A systematic study shows that only a certain subset of these operators can be the leading contribution to neutrino mass, if the Weinberg operator is forbidden by a discrete symmetry. We have discussed a specific implementation of a \( d = 7 \) operator that has potentially interesting phenomenological implications at the LHC, such as displaced vertices and lepton number violating signals. Those models have to be embedded in GUT multiplets in order to conserve gauge coupling unification. We have studied the decomposition of the \( d = 7 \) operator in an SU(5) framework, which is the only possible implementation of this operator consistent with the requirement of perturbativity. A particularly interesting realization of this model requires a softly broken symmetry. We have demonstrated that in this framework it is possible to obtain tri-bimaximal mixing with large corrections leading to non-zero \( \theta_{13} \). The soft breaking of the symmetry also leads to decays of additional heavy quarks, which are required in the GUT completion of this model. These decays are fast enough to be in agreement with cosmological bounds.

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