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\textbf{A B S T R A C T}

We consider extensions of the next-to-minimal supersymmetric model (NMSSM) in which the observed neutrino masses are described in terms of effective dimension six (or seven) rather than dimension five operators. All such operators respect the discrete symmetries of the model. The new particles associated with the double (or triple) seesaw mechanism can have sizable couplings to the known leptons, even with a TeV seesaw scale. In the latter case some of these new short-lived particles could be produced and detected at the LHC.

The next-to-minimal supersymmetric standard model (NMSSM) \cite{1} provides a well motivated extension of MSSM in which the $\mu$ problem of the latter is resolved through the introduction of a gauge singlet superfield $S$. This (NMSSM) extension has several phenomenological consequences. For instance, the upper bound on the lightest (CP-even) MSSM Higgs scalar can be increased from around 125 GeV to close to 140 GeV \cite{2}, and the little hierarchy problem encountered in the MSSM is ameliorated \cite{3,4}. New Higgs boson decay channels into CP-odd scalars appear \cite{4–6} which have an impact on Higgs boson searches, and there are new implications for dark matter physics \cite{7,8}.

In order to incorporate the observed solar and atmospheric neutrino oscillations \cite{9}, we propose in this Letter some extensions of the NMSSM based on the seesaw mechanism for providing neutrino masses. With the Large Hadron Collider (LHC) era about to unfold, we are especially interested in those seesaw extensions of the NMSSM which can be tested at the LHC. This leads us to consider dimension six and seven (rather than dimension five) operators for the generation of the observed neutrino masses and mixings. The presence of the $S$ VEV in the NMSSM turns out to be an important ingredient in implementing the double \cite{10} (or triple) seesaw mechanism. We consider several possibilities for renormalizable models at high energies, which may include SU(2) singlet, triplet and even additional doublet chiral superfields.

For the double or triple seesaw case, it is technically natural that even with a seesaw mass scale as low as 1 TeV or so, the new particles may have large couplings with the known leptons and Higgs doublets. This is generically not possible for type I or III seesaw. It is therefore an exciting possibility that some of the new particles we introduce could be produced at the LHC and detected through some distinctive decay signatures. We point out a way to experimentally distinguish the new particles involved in double seesaw from the ones in the conventional seesaw. We also consider the low energy implications of lepton-number conserving but lepton-flavor violating effective dimension six operators (in the Kähler potential).

We begin by recalling the basic structure of the NMSSM \cite{1}. We introduce a MSSM gauge singlet chiral superfield $S$ (with even $Z_2$ matter parity) through the following superpotential terms:

\begin{equation}
W \supset \lambda S H_u H_d + \frac{\kappa}{3} S^3.
\end{equation}

where $\lambda$ and $\kappa$ are dimensionless constants, and $H_u, H_d$ denote the MSSM Higgs doublets. A discrete $Z_3$ symmetry under which $S$ carries unit charge $\omega = e^{i2\pi/3}$ is introduced in order to eliminate from $W$ terms that are linear and quadratic in $S$, as well as the MSSM $\mu$ term. Note that $S$ could be assigned a $Z_3$ charge $\omega^3$, but this leads to the same dimension six and seven operators for neutrino masses. In order to decide on the $Z_3$ charges of the MSSM Higgs doublets, we require the presence in $W$ of Yukawa couplings at the renormalizable level. There are several possible $Z_3$ charge assignments for the matter superfields that are consistent with this requirement as displayed in Table 1. The $Z_3$ charge assignments in Table 1 for the quark superfields is not unique, but this will not be relevant for the discussion which follows.

We assume that the breaking of supersymmetry in the ‘hidden’ sector induces electroweak scale soft scalar mass terms consistent with $Z_3$ and $Z_2$:

\begin{equation}
V_{\text{soft}} = m_{H_u}^2 |H_u|^2 + m_{H_d}^2 |H_d|^2 + m_S^2 |S|^2 + (m_S S H_u H_d + m_S S^3 + \text{h.c.}).
\end{equation}
The scalar component of S acquires a non-zero VEV which generates the desired MSSM $\mu$ term. The radiative electroweak breaking scenario proceeds as in the MSSM case.

According to the charge assignments in Table 1, neutrino masses arise from effective dimension five, six or seven operators:

- **Cases la–lc**: $L H_u H_u \overline{L}$, $\frac{L H_u H_u S}{M_5}$, $\frac{L H_u H_u S^2}{M_6^2}$
- **Cases Ia–Ic**: $L H_u H_u \overline{L}$, $\frac{L H_u H_u S}{M_5}$, $\frac{L H_u H_u S^2}{M_6^2}$

where $M_{5,6,7}$ denote the appropriate seesaw mass scales.

### Case I

Cases la–lc correspond to the conventional dimension five neutrino operators. According to the three distinct ways to contract the SU(2) indices in Eq. (3), there are three kinds of seesaw mechanisms: type I seesaw mediated by the MSSM gauge singlet fermions, type II seesaw mediated by SU(2) triplet scalars with unit hypercharge, and type III seesaw mediated by SU(2) triplet fermions with zero hypercharge. In this Letter our focus will be mainly on dimension six and seven operators in Eqs. (4) and (5).

### Case II

Following electroweak symmetry breaking, the dimension six operator (Eq. (4)) induces light neutrino Majorana mass given by

$$m_{\nu} \sim \left( \frac{\nu_{2}^{2}}{M} \right) \times \left( \langle S \rangle / M \right).$$

where we set $\langle H_u \rangle \equiv v_u$. Compared to the conventional seesaw formula, we have an additional suppression factor $\langle S \rangle / M$ so for this case the upper bound on the seesaw mass scale is of order $10^5$ GeV, assuming all Yukawa couplings involved in the seesaw mechanism are of order unity. In practice, the seesaw scale can be much lower.

It is interesting and instructive to propose an explicit model which can generate these effective dimension six operators. We will show that the masses of some of the new particles we introduce can be within reach of the LHC.

Note that cases Ila–Iic give rise to the same dimension six operator given by Eq. (4). The heavy fields which generate dimension six operators for cases Ila–Iic will differ only in the choice of $Z_3$ charges. Thus, we will consider only case Ila, which is easily generalized for Iib and Iic.

As our first example on how to generate dimension six operators (Eq. (4)), we introduce the following new particles in the NMSSM (Fig. 1).

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### Case I

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As our first example on how to generate dimension six operators (Eq. (4)), we introduce the following new particles in the NMSSM (Fig. 1).
flavor-dependent corrections to the kinetic term of the left-handed neutrinos [14]. Taking again for simplicity $m_{ij} = M_0 \delta_{ij}$ and $(\lambda_N)_{ij} = \lambda_N \delta_{ij}$, the modified kinetic term of the left-handed neutrinos is found to be

\begin{equation}
L_{km} = i \sqrt{2} G_F (\delta_{ij} + \epsilon^k_{ij}) \nu^\mu \bar{\nu}_j \nu_{ik},
\end{equation}

where

\begin{equation}
\epsilon^k_{ij} = \frac{v^2}{M_0^2} (\nu^k \nu)_ij.
\end{equation}

In the presence of the flavor-dependent kinetic term, the relation between the mass $\bar{\nu}$ and flavor $\nu$ eigenstates of the light Majorana neutrinos is given by $\nu_i = (\mathcal{U}_{\nu \nu})_{ij} \bar{\nu}_j$, with $\mathcal{U}_{\nu \nu}$ the usual neutrino mixing matrix and

\begin{equation}
N_{ij} \simeq \delta_{ij} - \frac{1}{2} \epsilon^k_{ij},
\end{equation}

in the approximation $\epsilon^k_{ij} \ll 1$. Note that the matrix $\mathcal{U}$ is not unitary. This causes interesting modifications in both the charged and neutral currents (involving light neutrinos) in the SM:

\begin{equation}
J^{CC} \mu = \tilde{E}_{ij} \nu_i, \quad J^{NC} \mu = \frac{1}{2} \tilde{Y}_{ij} \lambda (\lambda' \nu')_{ij} \bar{\nu}_j.
\end{equation}

expressed in terms of the mass eigenstates. It turns out [15,16] that the elements of $|\lambda' \nu'|$ are somewhat severely constrained by the current experimental data on neutrino oscillations, $W$ and $Z$ boson decays, and flavor-violating decays of leptons: $Y_{ij} \lesssim 0.1$ for $M_0 \sim 1$ TeV. On the other hand, new signals of CP-violation related to this non-unitary leptonic mixing may be observed in future neutrino oscillation experiments [17].

An alternative way to generate the effective dimension six operator (Eq. (4)) is to replace the singlet superfields in Eq. (7) with two SU(2) triplets with zero-hypercharge:

\begin{equation}
\begin{array}{cccc}
SU(2) & U(1) & Z_3 & Z_2 \\
\Delta^c & 3 & 0 & \omega^2 \\
\Delta & 3 & 0 & \omega \\
\end{array}
\end{equation}

The superpotential in this case is given by

\begin{equation}
W \supset Y_{ij} (H_u \Delta^c_i L_j) + \frac{(\lambda_\Delta)_{ij}}{2} S \text{tr}[\Delta_i \Delta_j] + m_{ij} \text{tr}[\Delta_i^c \Delta_j].
\end{equation}

Integrating out the heavy triplets give rise to dimension six operators in the superpotential and the Kähler potential. Substituting the various VEVs, we obtain the light neutrino Majorana masses and flavor-violating kinetic terms. One difference from the previous (singlet) case is that dimension six operators are also induced for the charged leptons by integrating out the heavy charged fields in the SU(2) triplets. Again, in the mass basis, a non-unitary mixing matrix is induced, whose elements are constrained by the current experimental data [16]: $Y_{ij} \lesssim 0.01$–0.1 for $m_{ij} \sim 1$ TeV.

Yet another way for generating the dimension six operator is to introduce four additional SU(2) triplets with unit-hypercharge (see Fig. 3):

\begin{equation}
\begin{array}{cccc}
SU(2) & U(1) & Z_3 & Z_2 \\
\Delta^c & 3 & +1 & 1 \\
\Delta^c & 3 & -1 & 1 \\
\Delta & 3 & -1 & \omega \\
\tilde{\Delta} & 3 & +1 & \omega^2 \\
\end{array}
\end{equation}

The additional contributions to the NMSSM superpotential in this case contain the following terms

\begin{equation}
W \supset Y_{ij} (\bar{\Delta} \Delta^c L_j) + Y_{ih} (H_u \Delta^c H_u) + \lambda_N S \text{tr}[\Delta^c_i \tilde{\Delta}] + m_{ij} \text{tr}[\Delta^c_i \Delta_j] + m \text{tr}[\tilde{\Delta} \tilde{\Delta}].
\end{equation}

In this case, the effective Kähler potential after integrating out the heavy triplets is found to be of the form

\begin{equation}
\mathcal{K}_{\text{eff}} \sim Y_{ij} Y_{ij} \frac{L_i^1 L_j^1 L_k L_l}{m^2}.
\end{equation}

For the case of a non-unitary leptonic mixing a similar effective Kähler potential is generated. For a unitary leptonic mixing the effective Kähler potential takes the form (symbolically),

\begin{equation}
\mathcal{K}_{\text{eff}} \sim Y_{ij} Y_{ij} \frac{L_i^1 L_j^1 S}{m^2} + Y_{iA} Y_{ij} \frac{L_i^1 H_\mu^t H_\mu}{m_3^2},
\end{equation}

and the flavor-dependent kinetic terms are generated via the VEVs of $S$ and $H_u$. Comparing to the constraints on $Y_{ij}$ in the previous cases, we read the current experimental bounds as $Y_{ij} \lesssim 0.01$–0.1 for $m_{ih}, m_{A} \sim 1$ TeV.

### Case III

Cases IIIa–IIIc correspond to dimension seven operators for neutrino masses and mixings (see Eq. (5)). We note that this operator can be generated by integrating out the same heavy fields which we introduced for generating the dimension six operator. The main difference is in the $Z_2$ charge assignments. It is obvious that there are more possibilities to generate dimension seven operators for
neutrino masses compared to the dimension six case. We will provide one example of how to generate such an operator.

For case IIIa, we introduce the following new particles in the NMSSM spectrum (see Fig. 5),

\[
\begin{array}{c|ccc}
  \text{SU(2)} & U(1) & Z_3 & Z_2 \\
  \text{N}_j^c & 1 & 0 & \omega \sigma - \\
  \text{N}_j & 1 & 0 & \omega^2 \sigma - \\
  \text{N}_j^0 & 1 & 0 & 1 - \\
\end{array}
\]

(25)

where \(i, j\) denote the generation indices. To reproduce the neutrino oscillation data, we need to introduce at least two generations of \(N_j^c, N_j^0\) and \(N_j\).

The relevant part of the renormalizable superpotential involving only the new chiral superfields is given by

\[
W \supset Y_{ij} N_i^c (H_u L_j) + (\lambda_N N_i^c N_j^0 + m_{ij} N_j N_j^0 + \frac{1}{2} m_{ij}^0 N_i^0 N_j^0). \tag{26}
\]

For \(m_j^0\) and \(m_j\) larger than the electroweak scale, we integrate out the heavy \(N_j^c, N_j^0\) and \(N_j\) under the SUSY vacuum conditions. After eliminating the heavy fields, we arrive at the dimension seven operator of the form:

\[
W_{\text{eff}} = -\frac{1}{2} (H_u L_j)^T Y (m^1)^T (\lambda_N S)^T (m^0)^{-1} (\lambda_N S) m^{-1} Y (H_u L). \tag{27}
\]

For simplicity, we take \(m_{ij} = m_{ij}^0 = M_7 \delta_{ij}\) and \((\lambda_N N_i^c N_j^0 = \lambda_N \delta_{ij})\). Following the electroweak symmetry breaking, the neutrino Majorana mass matrix is generated:

\[
m_{ij} = \frac{(Y^T Y) v_u^2}{M_7^2} \times \frac{\lambda_N^2 \lambda_N (S^2)}{M_7^2}. \tag{28}
\]

We can see from the formula above that the upper bound for seesaw scale is \(M_7 \sim 10^6\) GeV, assuming all Yukawa coupling in Eq. (27) are \(\mathcal{O}(1)\). It is clear from Eq. (28) that by suitably adjusting the parameter \(\lambda_N\), the seesaw scale \(M_7\) can easily be lowered to the TeV range. This opens up the exciting possibility that some of the new particles in Eq. (26) can have \(\mathcal{O}(1)\) Yukawa couplings with the known matter fields.

If some of the new particles generating the seesaw mechanism have masses around 1 TeV, they could be produced in hadron colliders [18]. In type I seesaw, the heavy Majorana neutrino productions at hadron colliders via \(W\)-boson exchange and \(WW\)-fusion process [19–21] have been investigated. Signatures for the Majorana neutrinos could be observed through their lepton-number violating decays leading to like-sign dilepton production [20,21]. However, the production cross section is normally small, because the heavy neutrinos dominantly consist of the singlet Majorana neutrinos and the couplings between the heavy neutrinos and the \(W\)-boson are suppressed by a factor \(Y v_u/M_5\), the mixing between left-handed neutrinos and right-handed heavy neutrinos induced by the seesaw mechanism. Agreement with the neutrino oscillation data requires that, \(Y\) should at most be around \(10^{-5}\), and so the mixing is very small for \(M_5 \sim 1\) TeV. Some fine-tuning for the Dirac Yukawa matrix is necessary to keep the mixing angle as large as possible while reproducing the neutrino oscillation data [21].

The situation for type II or III seesaw is different, since the new particles involved in the seesaw mechanism have gauge couplings with the photon, \(W\) and 3-boson. Thus, these new particles could be produced at the LHC through processes mediated by these gauge bosons. The signature of such particles in type II seesaw [22] and in type III seesaw [23] have been investigated in detail. For example, the doubly-charged scalar in type II seesaw, once produced, may provide a clean signature through its decay into a pair of same sign charged leptons [22]. The pair production of the singly-charged fermions in type III seesaw and their (lepton-number and/or lepton-flavor violating) decays into leptons and \(W\), \(Z\)-boson or Higgs boson could be discovered at the LHC. Note that if the seesaw scale is around 1 TeV, the Dirac Yukawa couplings are very small, say, \(Y = \mathcal{O}(10^{-5})\) to provide a light neutrino mass \(m_\nu = \mathcal{O}(1\, \text{eV})\). This fact implies that the new particle production is accompanied by an extra signature [23]: The lifetime of the produced particles is long enough for the decay vertices to be detectably displaced from the primary production vertex.

The new particles included in the conventional seesaw models also appear in models with the double (or triple) seesaw mechanism. The collider phenomenology for these particles is analogous to the one in the conventional seesaw models. However, there is a crucial difference arising from the structure of the double (or triple) seesaw mechanism. As we have already noted, in the double (or triple) seesaw mechanism we can reproduce the light neutrino Majorana masses while keeping the Dirac Yukawa couplings as large as possible, so long as the couplings among the new particles are suitably small. Therefore, the production cross section of the heavy neutrinos in type I seesaw can be sizable. If the Dirac Yukawa couplings are much larger than the values expected in the conventional seesaw models with a TeV seesaw scale, the lifetimes of the new particles can be relatively short. For a low seesaw scale this could be a distinguishing feature between the conventional seesaw and the double (or triple) seesaw models. In addition, we have seen that in models with double (or triple) seesaw, new SU(2) doublets, which do not appear in the conventional seesaw models, can be present. The phenomenology of these new particles would be worth investigating. For example, the charged scalar in the doublet superfield \((H_u^c)\), once produced, can decay into charged-leptons and the fermionic component of the singlet superfield \(S\).

Recently, it has been pointed out [8] that if the right-handed neutrinos have couplings with the singlet \(S\) in the NMSSM, the lightest right-handed sneutrino can be a viable cold dark matter candidate through its coupling with the Higgs bosons. Our double (or triple) seesaw mechanism with right-handed neutrinos shares the same structure for the right-handed sneutrinos, so that this scenario can also work in our models. Finally, let us recall that the presence of low scale seesaw can alter the predictions for the SM Higgs boson mass [24]. In the NMSSM case, the coupling \(Y H_i (H_u \Delta H_u)\) in Eq. (20) will generate a tree level contribution to the lightest CP-even Higgs boson mass [25]:

\[
m_h^2 = (m_H^2)_{\text{NMSSM}} + 4 Y^2 v_u^2 \sin^2 \beta. \tag{29}
\]

where \(\tan \beta = v_u/v_d\) and \((m_H^2)_{\text{NMSSM}}\) denotes the standard NMSSM contribution. There is a constraint on the triplet Higgs VEV arising from a global fit of electroweak data [26]:

\[
v_A \lesssim 2\, \text{GeV}. \tag{30}
\]

If one assumes that the triplet Higgs mass is a few TeV, the trilinear coupling \(Y_H\) can be order unity. In this case the tree level contribution can help to make the Higgs mass as heavy as 200 GeV (for \(\tan \beta > 10\)). For a more precise estimate, loop correction arising from the Higgs triplet should be included.
The philosophy behind our construction of an extended NMSSM which yields dimension five, six or seven seesaw operators is identical to the one employed in the standard NMSSM construction. Namely, we require that all renormalizable superpotential couplings respect the $Z_3$ symmetry. The higher dimensional seesaw operators arise from integrating out the new heavy states. As far as the $Z_3$ domain wall problem is concerned we do not have anything new to add to the discussion which appears in the literature. Thus, we are assuming that the solution suggested by C. Panagiotakopoulos and K. Tamvakis [27] is also applicable to our case, perhaps with suitable modification. That is, suitable higher dimensional operators which we have discussed in our Letter.

In summary, we have proposed extensions of the NMSSM particle spectrum such that the observed neutrino masses are described in terms of effective dimension six or seven (instead of dimension five) operators. In addition to the models corresponding to type I, II and III seesaw, we introduce more exotic possibilities with SU(2) doublet superfields which do not appear in the conventional seesaw models. The new heavy states responsible for the seesaw mechanism can have sizable couplings to the known leptons and may be detected at the LHC. The low energy implications of lepton flavor violating dimension six operators in the Kähler potential are briefly discussed.

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