Neutrino Mass Seesaw at the Weak Scale, the Baryon Asymmetry, and the LHC

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We consider theories where the Standard Model (SM) neutrinos acquire masses through the seesaw mechanism at the weak scale. We show that in such a scenario, the requirement that any pre-existing baryon asymmetry, regardless of its origin, not be washed out leads to correlations between the pattern of SM neutrino masses and the spectrum of new particles at the weak scale, leading to definite predictions for the LHC. For type I seesaw models with a TeV scale $Z'$ coupled to SM neutrinos, we find that for a normal neutrino mass hierarchy, at least one of the right-handed neutrinos must be `electrophobic', decaying with a strong preference into final states with muons and tauons rather than electrons. For inverted or quasi-degenerate mass patterns, on the other hand, we find upper bounds on the mass of at least one right-handed neutrino. In particular, for an inverted mass hierarchy, this bound is 1 TeV, while the corresponding upper limit in the quasi-degenerate case is 300 GeV. Similar results hold in type III seesaw models, albeit with somewhat more stringent bounds. For the Type II seesaw case with a weak scale $SU(2)$ triplet Higgs, we again find that an interesting range of Higgs triplet masses is disallowed by these considerations.

\section{Introduction}

Neutrino masses constitute concrete evidence for the existence of physics beyond the Standard Model (SM). However, much about the neutrino sector remains to be understood. It is not known whether neutrinos are Dirac or Majorana, or whether the pattern of neutrino masses is hierarchical, inverse hierarchical or quasi-degenerate. Further, the dynamics by which the neutrinos acquire their masses, and the scale at which this occurs, are also not known.

One interesting possibility is that the SM neutrinos acquire their masses at or close to the weak scale through the seesaw mechanism \cite{See}. In such a scenario the dynamics underlying neutrino mass generation might be accessible to the LHC \cite{LHC}. This could happen in several different ways. For example, if the SM neutrinos are charged under a new $U(1)$ gauge symmetry that is broken only close to the weak scale, they are prevented from acquiring mass at higher scales. In these theories a Type I seesaw at the weak scale is perhaps the simplest possibility for neutrino mass generation, with some of the additional particles required for cancellation of the $[U(1)]^3$ anomalies playing the role of right-handed (RH) neutrinos. Such a scenario is promising for the LHC because the RH neutrinos could be pair-produced through decays of the $Z'$ gauge boson, and are in general straightforward to detect through their lepton-number-violating decaying states. Other classes of seesaw models that are exciting for the LHC are Type II models with an $SU(2)$ triplet Higgs at the weak scale \cite{TypeII}, and Type III models with $SU(2)$ triplet RH neutrinos at the weak scale \cite{TypeIII}.

In this letter we show that in theories where the Standard Model (SM) neutrinos acquire Majorana masses through the seesaw mechanism at the weak scale, the requirement that any pre-existing baryon asymmetry is not almost entirely washed out can be used to correlate the pattern of SM neutrino masses to the spectrum of new particles at the weak scale. Then, unless the baryon asymmetry is generated at or below the weak scale, or the pre-existing baryon asymmetry is extremely large, this leads to definite predictions at the LHC for each of the three classes of seesaw models described above.

- In the case of Type I seesaw models with a weak-scale $Z'$, if the SM neutrinos are hierarchical, at least one of the RH neutrinos must be `electrophobic', decaying with a strong preference into final states with muons and tauons rather than electrons. On the other hand, if the SM neutrinos exhibit an inverted (quasi-degenerate) pattern of masses, at least one of the RH neutrinos must be lighter than 1 TeV (300 GeV).
- For Type III seesaw models, the corresponding upper bound for $SU(2)$ triplet RH neutrinos is 300 GeV for inverted neutrino mass pattern and 170 GeV for the quasi-degenerate case.
- In the case of Type II seesaw models, a significant portion of the triplet VEV and mass $(\nu_{\Delta}, M_{\Delta})$ parameter space accessible at the LHC is disfavored.

It should be noted that our objective is different from that of previous works whose purpose was to explain the matter-antimatter asymmetry thanks to leptogenesis (for a recent review, see \cite{Leptogenesis}). Leptogenesis in the Type I case has been extensively studied, and it was found that there exists a lower bound on the scale of leptogenesis around $10^9$ GeV in the limit of hierarchical RH neutrinos \cite{TypeI,Leptogenesis}. In the Type II case, \cite{TypeII} found that there is also a lower bound which points to a very high scale. Similarly, the Type III was worked out in \cite{TypeIII}, where a similar bound was obtained. Let us stress here that these bounds are in no way inconsistent with our approach, which is
The baryon asymmetry originated. existing baryon asymmetry, regardless of how the baryon can, in combination with sphalerons [12], erase any pre-violation at low energies. If lepton-number-violating will generally lead to some amount of lepton number violation at low energies. If lepton-number-violating processes are in equilibrium at the weak scale, they can, in combination with sphalerons [12], erase any pre-existing baryon asymmetry, regardless of how the baryon asymmetry originated.

As a concrete example, consider a toy Type I model with a single SM lepton doublet $L$ that acquires a Majorana mass after coupling through the Higgs $H$ to a single RH neutrino $N$ that has a weak-scale mass $M$. The relevant part of the Lagrangian is

$$\lambda LHN + MN^2 + \text{h.c.}$$  \hspace{1cm} (1)

At temperatures $T \geq M$ the rate for lepton-number-violating decays and inverse decays of $N$ is given by

$$\Gamma \sim \lambda^2 \frac{M^2}{T} \sim m_{\nu} \frac{M^3}{v^2 T}$$  \hspace{1cm} (2)

where $v = \langle H \rangle$ is the SM Higgs VEV and $m_{\nu} \sim \lambda^2 v^2 / M$ is the SM neutrino mass. For this process to be in equilibrium, this rate must be larger than the expansion rate $H \sim T^2 / M_{\text{Pl}}$, where $M_{\text{Pl}}$ is the Planck scale. Lepton-number-violating decays will be in equilibrium for $T \sim M$ if $m_{\nu} \gtrsim v^2 / M_{\text{Pl}}$. For $m_{\nu}$ of order 0.05 eV, the atmospheric scale, this condition is satisfied in our toy model, and any pre-existing baryon asymmetry will endure washout. However, whether or not the baryon asymmetry is erased in realistic seesaw models is a far more detailed question. There are several reasons for this.

- Since the medium distinguishes all three lepton flavors $L_e$, $L_\mu$, and $L_\tau$ at the weak scale, and since sphalerons conserve $B/3 - L_\alpha$, ($\alpha = e, \mu, \tau$) lepton flavor violating processes must be in equilibrium for each lepton flavor in order for the baryon asymmetry to be efficiently erased. This requires taking into account the full spectrum of neutrino masses and mixings.

- The extent to which the baryon asymmetry is erased depends on the period of time that the lepton-flavor-violating processes, most importantly inverse decays, are in equilibrium. For $m_{\nu}$ of order the atmospheric scale, these processes are only in equilibrium for a short time. It then requires a careful analysis to determine how much of the original baryon asymmetry survives. The relevant Boltzmann equation describing the washout process by the seesaw mediator(s) of mass $M$ can be written in the general form

$$\frac{dY_{B/3-L_\alpha}}{dz} = -W_\alpha(z)Y_{B/3-L_\alpha},$$  \hspace{1cm} (3)

where $z = M/T$, $Y_{B/3-L_\alpha}$ is the $B/3 - L_\alpha$ number density over the entropy density, and $W_\alpha$ is a generic washout term which will have to be specified in each particular case. This formulation of the washout term is consistent with leptogenesis analyses in the Type I and III cases [3, 11] by setting the $CP$ asymmetry to zero. In the Type II case [14], the situation is slightly more complicated since the scalar triplets carry non-zero hypercharge. As we will show explicitly in Section V in that case, the evolution of the scalar triplet asymmetry must be tracked as well. However, we will find that the more complicated system of equations reduces to this simple form in a certain limit.

The solution to Eq. (3) can be easily obtained

$$Y_{B/3-L_\alpha}(z) = Y_{B/3-L_\alpha}^{\text{in}} \exp \left[-\int_{z_{\text{in}}}^{z} dz' W_\alpha(z') \right],$$  \hspace{1cm} (4)

where $Y_{B/3-L_\alpha}^{\text{in}}$ stands for any pre-existing asymmetry at $z_{\text{in}} \ll 0.1$. One can therefore see that any pre-existing asymmetry will be exponentially washed out provided the integral is large.

- The baryon asymmetry, having been erased, can in principle be regenerated by late out-of-equilibrium decays of RH neutrinos. However, if the masses of the RH neutrinos are of order the weak scale as in the case of interest, the relevant Yukawa couplings are generally much too small to allow a large-enough asymmetry to be generated. However, an exception to this rule occurs if the the RH neutrinos are extremely degenerate, at the level of one part in $10^{10}$, in which case resonant leptogenesis can occur [13]. In what follows we will assume that the RH neutrinos are not sufficiently degenerate to allow this possibility, leaving this very special case for future work.

In the subsequent sections we analyse each of the three classes of seesaw models in turn, and determine the precise conditions under which the baryon asymmetry is erased, and the consequences for LHC phenomenology. We shall place limits on the spectrum of particles at the weak scale by requiring that the baryon asymmetry not be washed out by a factor greater than $10^9$. It is important to note that these bounds can be evaded
the baryon asymmetry is generated below the weak scale, or alternatively, if the primordial baryon asymmetry generated at high scales is extremely large, of order $10^{-3}$ or more. However, such a large primordial asymmetry cannot be generated by the familiar out-of-equilibrium decays of a heavy field, but instead requires a very efficient mechanism of baryogenesis, such as the one proposed by Affleck and Dine \[17\]. If the LHC were to discover that these bounds are violated, it would shed new light on the mechanism that generates the baryon asymmetry of the Universe.

II. THE TYPE I SEESAW

For the Type I seesaw case, we assume the presence of three RH neutrinos $N_i$, $i = 1, 2, 3$ with masses $M_i$ around the weak scale. The washout parameter introduced in Eq. \[8\] is given in this case by

$$W_{iα}^α(z_i) = \sum_j W_{iα}^{ID}(z_i),$$

where $z_i = z M_i/M_1$. The washout is dominated by inverse decays $W^{ID}$, which can be conveniently expressed as

$$W_{iα}^{ID} = \frac{1}{4} K_{iα} K_1(z_i) z_i^3,$$

where $K_1(z)$ is the modified Bessel function, and

$$K_{iα} = \frac{\Gamma_D(N_i \to ℓ_αH + ℓ_αH^\dagger)}{H(z_i = 1)} = \frac{|λ_{iα}|^2 v^2}{M_i m_*},$$

with $m_* = 1.08 \times 10^{-3}$ eV.

We show in Fig. 1 how the function $W_{iα}(z_i)$ behaves for different values of $K_{iα}$. Note that an identical figure would be obtained in the Type II and III cases, where only the definition of $K_{iα}$ is different. It is useful to derive the maximum washout possible, which from Fig. \[1\] is obtained for $M \gtrsim 10^7$. In this limit ($z \gg 1$), the integral in Eq. \[6\] can be solved analytically \[15\]:

$$Y_{B/3−L_α}(∞) = Y_{B/3−L_α}^{in} \exp \left[ \frac{3π}{8} \sum_i K_{iα} \right].$$

However, more precisely, the upper bound on the integral in Eq. \[6\] should be set by the decoupling temperature of the sphalerons. For $M_H = 120$ GeV, it is given by $T_{dec} \simeq 130$ GeV \[19\]. At fixed temperature $T = T_{dec}$, it can be seen from Fig. \[1\] that the area below the curve $W(z)$ increases as $M$ increases, so that the washout increases as well [see Eq. \[5\]]. Therefore, by imposing that the washout does not exceed a factor of $10^6$, we will obtain upper bounds on $M$.

The question then is: how to know the magnitude of $K_{iα}$. This can be answered using the parametrization of the Yukawa coupling matrix \[17\]

$$λ_{iα} = \left( U \sqrt{D_mΩ} \sqrt{D_M} \right)_{iα}/v,$$

where $U$ is the PMNS mixing matrix, $D_m$ and $D_M$ are diagonal matrices for the masses of light and heavy neutrinos, respectively, and $Ω$ is a complex orthogonal matrix. Using this parametrization, one can express the decay parameters in Eq. \[8\] in the following way:

$$K_{iα} = \frac{1}{m_*} \left| \sum_j \sqrt{m_j} U_{ij} Ω_{ji} \right|^2.$$

It should be noted that $K_i = \sum_α K_{iα}$ obey the following inequality:

$$\sum_i K_i \gtrsim \frac{1}{m_*} \left| \sum_α Ω_{ji} \right|^2.$$

Without flavor effects, the washout implied would be huge. With flavor effects the washout is reduced, but one has to check to what extent.

In the case of normal hierarchy and $m_1 \ll m_{sol}$, due to the small $U_{e3}$ entry in the PMNS matrix, the washout in flavor $e$ is typically suppressed, and it can be as low as $\sum_i K_{iα} \sim 2$, implying little washout. Clearly, the washout in the other two flavors is very effective. We show below that in this case at least one RH neutrino will decay electrophobically. Let us explain how this comes about by first writing the branching ratio for the decay of $N_i$ into each flavor $α$:

$$B_{iα} = \frac{\left| \sum_j \sqrt{m_j U_{ij} Ω_{ji}} \right|^2}{\sum_j |m_j Ω_{ji}|^2}.$$

In the case of normal hierarchy with $m_1 \ll m_{sol}$, for flavor $e$ the largest mass eigenstate $m_3 = m_{atm}$ is coupled to the small entry $U_{e3} < 0.2$ (3σ). One can then envisage three
We obtain that 99% of the points lie above $\sum_i K_{i\alpha} = 12$.

situations: $|\Omega_3| \ll |\Omega_2|$, $|\Omega_3| \sim |\Omega_2|$, or $|\Omega_3| \gg |\Omega_2|$. In the first case, the decay will be roughly 1/3 in each flavor, but the total rate will be suppressed since it is driven by the solar scale. In the second case, both terms are comparable, implying that $N_i$ will decay typically only 8% into $e$. In the third case, one obtains $B_{\alpha i} \simeq |U_{ei}|^2 < 4\%$. There is however a loophole in the first case. From the orthogonality of $\Omega$, $\sum_i \Omega_{ij}^2 = 1$, it is not possible to have $|\Omega_3| < |\Omega_2|$ for all $i$. Hence, for at least one RH neutrino, the first case is excluded and the decay will be electrophobic.

For an inverted hierarchy and $m_3 \ll m_{\text{sol}}$, the situation is different because both $m_1$ and $m_2$ are roughly at the atmospheric scale, implying typically a larger washout than in the previous case. Quantitatively, varying $0 \leq |\Omega_{ij}| \leq 1$ as well as the unknown phases in the PMNS phases, we find that 99% of the points satisfy $\sum_i K_{i\alpha} \geq 12$. In Fig. 2 we show the scatter plot from which this lower bound was extracted. This means that the baryon asymmetry will be washed out by more than a factor of $10^6$ when $M_1 > 1$ TeV for 99% of the parameter space. Thus we consider conservatively 1 TeV to be the upper bound on the mass of the lightest RH neutrino. Note that this result also holds when $|\Omega_{ij}|$ is allowed to take larger values, because the washout increases as $K \propto |\Omega|^2$.

For a quasi-degenerate spectrum, $m_1 \simeq m_2 \simeq m_3 \simeq 0.1$ eV, the washout is expected to be even larger. It turns out that $\sum_i K_{i\alpha} \geq 40$, implying an upper bound $M_1 < 300$ GeV, a constraint stronger than the inverted hierarchy case. We present in Fig. 3 a plot of the washout factor in Eq. (4) versus the RH neutrino mass. This figure shows explicitly how the upper bound of about 300 GeV was obtained.

A. Neutrino textures

We now comment on the possibility of evading the bounds derived above in the case of inverted hierarchy or quasi-degenerate spectrum by a suitable texture in the Yukawa coupling matrix for neutrinos. The question is whether the points in the parameter space which we did not consider (e.g. points for which $\sum_i K_{i\alpha} \ll 12$ in Fig. 3) require fine-tuning of model parameters to get desired neutrino masses and mixings or not. Our conclusion will be that the structure needed in order to avoid the washout is not compatible with an inverted hierarchy or a quasi-degenerate spectrum, except with a very careful tuning of the parameters.

The most minimal condition for evading our bounds in the case of inverted or quasi-degenerate spectra is to reduce the washout in a particular flavor, as given below:

$$K = \left( \begin{array}{ccc} \epsilon_1^2 & \epsilon_2^2 & \epsilon_3^2 \\ a_1^2 & a_2^2 & a_3^2 \\ b_1^2 & b_2^2 & b_3^2 \end{array} \right), \left( \begin{array}{ccc} \epsilon_1^2 & \epsilon_2^2 & \epsilon_3^2 \\ a_1^2 & a_2^2 & a_3^2 \\ b_1^2 & b_2^2 & b_3^2 \end{array} \right), \left( \begin{array}{ccc} \epsilon_1^2 & \epsilon_2^2 & \epsilon_3^2 \\ a_1^2 & a_2^2 & a_3^2 \\ b_1^2 & b_2^2 & b_3^2 \end{array} \right),$$

where we impose that $\sum_i |\epsilon_i|^2 \ll \sum_i |a_i|^2$, $\sum_i |b_i|^2$. Let us now calculate the form of the light neutrino mass matrices from the above $K$ matrix structure using Type I seesaw and diagonal RH neutrino mass matrix. We obtain the following three different forms:

$$m_\nu = v^2 \left( \begin{array}{cc} \lambda \frac{\lambda}{M} \lambda^T \end{array} \right),$$

$$m_\nu = \left( \begin{array}{ccc} O(\epsilon^2) & O(\epsilon) & O(\epsilon) \\ O(\epsilon) & A & C \\ O(\epsilon) & C & B \end{array} \right),$$

$$m_\nu = \left( \begin{array}{ccc} A & O(\epsilon) & C \\ O(\epsilon) & O(c^2) & O(\epsilon) \\ C & O(\epsilon) & B \end{array} \right).$$
Due to the condition \(\sum |\epsilon_i|^2 \ll \sum |a_i|^2\), \(\sum |b_i|^2\), it can be immediately seen that one eigenvalue will necessarily be suppressed, and therefore it not consistent with a quasi-degenerate spectrum.

For the case of inverted hierarchy, it can be shown that the condition to have one zero eigenvalue and two degenerate eigenvalues, i.e. \(m_3 = 0\) and \(m_1 = m_2\), is given by \(|A| = |B|\) and \(|C| = 0\). Allowing for a small perturbation, we have a matrix of the following form [for illustration we show only the first matrix in Eq. (14)]:

\[
m_\nu = m_* \begin{pmatrix}
O(\epsilon^2) & O(\epsilon) & O(\epsilon) \\
O(\epsilon) & O(\epsilon) & O(\epsilon) \\
O(\epsilon) & O(\epsilon) & A + O(\epsilon)
\end{pmatrix}
\]

(18)

Note that the potential phase difference between \(A\) and \(B\) can be absorbed by making the transformation \(\nu_1 \to e^{i\phi} \nu_1\). Diagonalizing a mass matrix of this form, it is easy to show that one mixing angle, \(\theta_{23}\), can be large, whereas the other two, \(\theta_{12}\) and \(\theta_{13}\) have to be small because of the smallness of \(\epsilon\). This is at odds with the observed large solar mixing angle, which, though not maximal, is large.

1. Examples of Neutrino Mass Models

It seems from the previous discussion that the observed neutrino masses and mixings constrain the possible patterns of the \(K\) matrix in such a way that a small washout in one flavor is not expected for inverted and quasi-degenerate spectra. In order to further illustrate this point, we give two examples of known neutrino models, which lead to different neutrino mass hierarchies.

a. \(L_e - L_\mu - L_\tau\) symmetry case: The well-known symmetry \(L_e - L_\mu - L_\tau\) predicts an inverted spectrum. One way to implement this symmetry is to use \(3 \times 2\) seesaw with the following Yukawa coupling matrix:

\[
\lambda = \begin{pmatrix}
a_1 & O(\epsilon) \\
O(\epsilon) & a_2 \\
O(\epsilon) & a_3
\end{pmatrix}.
\]

(19)

The \(O(\epsilon)\) terms denote small breaking of this symmetry. The \(L_e - L_\mu - L_\tau\) charges of \(N_{e,\mu}\) are \((+1, -1)\) respectively. The RH neutrino mass matrix in the flavor basis is given by

\[
\begin{pmatrix}
0 & M \\
M & 0
\end{pmatrix}.
\]

(20)

In the basis where the RH neutrino masses are diagonal, it can be immediately seen that the \(K\) matrix,

\[
K = \frac{v^2}{\sqrt{2} M m_*} \begin{pmatrix}
|a_1|^2 & |a_1|^2 \\
|a_2|^2 & |a_2|^2 \\
|a_3|^2 & |a_3|^2
\end{pmatrix},
\]

(21)

does not allow the washout to be avoided in any flavor.

One can also have examples with \(3 \times 3\) seesaw with obvious assignments under \(L_e - L_\mu - L_\tau\) symmetry for the lepton doublet and singlet fields. In the strict symmetry limit, the RH neutrino mass matrix is singular. One can therefore consider a softly broken \(L_e - L_\mu - L_\tau\) symmetry with a RH neutrino mass matrix of the following form:

\[
\begin{pmatrix}
0 & M_1 & M_2 \\
M_1 & 0 & 0 \\
M_2 & 0 & M_3
\end{pmatrix}.
\]

(22)

where the \(M_3\) term is not expected for inverted and quasi-degenerate spectra.

b. Quasi-degenerate \(A_4\) model: The above conclusion holds also in models that predict quasi-degenerate neutrinos. As an example, consider a model based on discrete non-Abelian \(A_4\) symmetry

\[
\lambda = \frac{1}{\sqrt{3}} \begin{pmatrix}
1 & 1 & 1 \\
1 & \omega & \omega^2 \\
1 & \omega^2 & \omega
\end{pmatrix},
\]

(26)

where \(\omega\) is a phase. Since the moduli of all Yukawa couplings are of the same order, any spectrum of RH neutrino masses will lead to a strong washout in all flavors.

In summary, realistic neutrino mass models do not lead to dramatic flavor effects. This is certainly true for inverted and quasi-degenerate neutrino mass spectra. Only for normal hierarchy, a reduced washout by an order of magnitude in the electron flavor can be naturally achieved. This is the reason why we did not obtain any constraint in that case.
III. TYPE III SEESAW

The Type III seesaw has the same parameters as the Type I, namely a neutrino Yukawa coupling matrix $\lambda$ and a Majorana mass matrix $M$ which can be taken to be diagonal. The essential difference is that the seesaw mediators are fermionic $SU(2)$ triplets, with one neutral and two charged components. Leptogenesis in this model was studied in [11]. In particular, the washout parameter in Eq. (3) is given by

$$W_{\alpha}^{III}(z) = \frac{1}{4} \sum_i K_{\alpha i}^{III} K_1(z_i) \xi_i^3,$$

where $K_{\alpha i}^{III} = 3 K_{i\alpha}$ because the three components of the triplet contribute to the washout.

We now study the constraints on the spectrum of the triplets from washout arguments. In the case of normal hierarchy and $m_1 \ll m_{\text{sol}}$, a pre-existing asymmetry in flavor $e$ will again partially escape the washout, although not as much as in the Type I case: a maximum washout factor of $10^4$. Second, for the same reason as in the Type I case, at least one of the triplets will decay electrophobically.

For an inverted hierarchy and $m_3 \ll m_{\text{sol}}$, we have $\sum_i K_{\alpha i}^{III} \geq 36$. We then obtain the conservative upper bound $M_1 < 300$ GeV, substantially more stringent than the corresponding type I case.

For a quasi-degenerate spectrum, $m_1 \approx m_2 \approx m_3 \approx 0.1$ eV, we obtain that $\sum_i K_{\alpha i}^{III} \geq 120$. This leads to the upper bound $M_1 < 170$ GeV.

It is worth pointing out that the LHC has the capacity to observe the triplet fermions up to 1 TeV in five years of operation [20]. A considerable portion of this allowed range would therefore be in conflict with high-scale baryogenesis.

IV. TYPE II SEESAW

In the Type II seesaw, neutrino masses are generated by the VEV of the neutral component of scalar $SU(2)$ triplet $\Delta$. For our purposes, the relevant part of the Lagrangian is

$$-h_{\alpha \beta} \ell_{L \alpha}^T C \sigma_2 \Delta L_{\beta} - \mu H^T i \sigma_2 \Delta^H + h.c.,$$

where $\alpha, \beta = e, \mu, \tau$, and the neutrino mass matrix is given by

$$(m_\nu)_{\alpha \beta} = h_{\alpha \beta} v_\Delta = h_{\alpha \beta} \frac{\mu v^2}{M_\Delta^2}. $$

It should be noted that only one triplet is able to generate three active neutrino masses, contrary to the Type I and III cases.

We are interested in the case where $\Delta$ mass in the 100 GeV to TeV range so that it is accessible at LHC [21]. We vary the parameter $\mu$ so that the triplet VEV $v_\Delta = \mu v^2 / M_\Delta^2$ also correspondingly changes. Note that electroweak $\rho$-parameter constraint implies $v_\Delta \leq 1$ GeV. We will always stay far below this value. Depending on the magnitude of $\mu$, the decay of the $\Delta^{++}$ will either be to a Higgs pair $\Delta \rightarrow H^+ H^+$ or to a like-sign dilepton pair $\Delta \rightarrow \ell \ell$. Which lepton pair will of course depend on the neutrino mass hierarchy; specifically, we expect the dominant channels to be $\tau \tau$, $\mu \tau$ and $\mu \mu$ for normal hierarchy and $e \mu$ and $e \tau$ or $ee$ for inverted hierarchy.

Due to the presence of the scalar triplet, which carries a non-zero hypercharge, the washout is here evaluated solving a set of Boltzmann equations, as presented in [10]. More precisely, coupled equations for the evolution of the asymmetries in the Higgs field, in the scalar triplet, as well as in each lepton flavor must be solved. Including flavor effects and the most relevant spectator processes, i.e. the Yukawa interactions *, and imposing hypercharge conservation, which allows one to remove one equation, we obtain the following system of equations:

$$\frac{dY_{e\alpha}}{dz} = -2D_\Delta B_{L\alpha}(Y_{\alpha} + 3z^2K_2(z)Y_{e\alpha}),$$

$$\frac{d\Delta}{dz} = -2D_\Delta \left[ Y_{\alpha} + 3z^2K_2(z) \left( \sum_\alpha B_{L\alpha} Y_{e\alpha} \right) \right] -3z^2K_2(z)B_H \frac{2}{3N_f} \left( -2Y_{\Delta} + \frac{8}{3} \sum_\alpha Y_{e\alpha} \right),$$

where $N_f$ is the number of quark generations, and where we defined

$$D_\Delta = \frac{\Gamma_\Delta K_1(z)/K_2(z)}{H z}$$

with the total triplet decay rate $\Gamma_\Delta = \frac{M_\Delta^2}{4\pi} \left( \sum_i m_i^2 / v_\Delta^2 + v_\Delta^2 M_\Delta^2 / v^4 \right)$. The branching ratios in each flavor $\alpha$ are given by

$$B_{L\alpha} = \frac{\sum_k m_k^2 |V_{\alpha k}|^2}{\sum_k m_k^2 + v_\Delta^2 M_\Delta^2 / v^4},$$

where $V$ is the PMNS matrix without Majorana phases, and $\sum_\alpha B_{L\alpha} = B_L = 1 - B_H$.

We solve Eqs. (32) and (31) with initial asymmetries in the lepton fields and Higgs fields so as to satisfy hypercharge neutrality!, and we obtain the region of parameters $(v_\Delta, M_\Delta)$ where the washout is greater than a factor of $10^6$. The result is shown in Fig. 4 for the different neutrino mass spectra. It can be seen that a substantial portion of the parameter space is disfavored by our washout argument in the case of inverted and

* We did not include sphaleron effects, which only induce small corrections.

† This is a conservative choice of initial conditions. An initial asymmetry in the scalar triplet fields would lead to more stringent limits.
Yukawa interactions, which redistribute the Higgs asymmetry among all SM particles, and which are at the origin of the factor $\frac{\Delta}{m^2}$ in Eq. (31), are absolutely essential to get the correct result. Using strictly the equations proposed in [10] adding only flavor effects, and therefore neglecting these Yukawa interactions, leads to an overestimation of the washout by up to 12 orders of magnitude! The bounds presented in Fig. 4 would then become more stringent in the region $v_\Delta \lesssim 10^{-4}$ by up to a factor of 2.

V. CONCLUSION

We studied in detail the consequences of weak-scale seesaw mechanisms on any pre-existing baryon asymmetry. If it is not to be efficiently washed out, we found correlations between the pattern of neutrino masses and the spectrum of new particles at the weak scale. For type I seesaw models with a TeV scale $Z'$ coupled to SM neutrinos, we found that for a normal neutrino mass hierarchy, at least one of the RH neutrinos must be ‘electrophobic’, decaying with a strong preference into final states with muons and taus rather than electrons. For inverted or quasi-degenerate mass patterns, on the other hand, we found upper bounds on the mass of at least one RH neutrino. In particular, for an inverted mass hierarchy, this bound is 1 TeV, while the corresponding upper limit in the quasi-degenerate case is 300 GeV. Similar results hold in type III seesaw models, albeit with somewhat more stringent bounds. For the Type II seesaw case with a weak scale $SU(2)$ triplet Higgs, we again found that an interesting range of Higgs triplet masses is disallowed by these considerations.

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