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Seesaw Theories at LHC and Warm Dark Matter

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Abstract. Explaining the origin of neutrino masses clearly requires new physics beyond the Standard Model. I focus on the Seesaw paradigm and discuss a few simplest extensions of the SM that give Majorana masses to the active neutrinos. If realized at TeV scale, seesaw theories could manifest themselves in lepton number violating signatures at both low-energy processes and high-energy collider experiments. I summarize the constraints on the seesaw scales using the current LHC data. The left-right symmetric model connects the seesaw mechanism with the origin of parity symmetry breaking, and provides a unified framework for the simplest seesaw types. With new right-handed charged-current interactions, a TeV such model offers a plethora of new particles and exotic signatures at the LHC, and also accommodates a dark matter candidate, the lightest right-handed neutrino. A challenging question is the dark matter relic density which is typically over-produced in the early universe. The late decays of two heavier right-handed neutrinos can produce entropy and dilute the dark matter number. The key observation for this picture to work is the interplay between the freeze temperature of TeV right-handed gauge interaction and the QCD phase transition. The resulting dark matter mass is predicted to be around keV which makes the left-right model also a theory of warm dark matter. I will also comment on the fate of cosmic baryon asymmetry in this scenario.

Keywords: Seesaw Mechanism, Left-Right Symmetry, Dark Matter, QCD Phase Transition, Entropy Production
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INTRODUCTION

The existence of non-zero neutrino masses observed in neutrino oscillations is an unambiguous evidence for physics beyond the Standard Model (SM). One possibility is neutrinos having Majorana masses [1]. This is based on the observation that a Weyl spinor is in the irreducible fermion representation of Lorentz transformation and the Majorana mass term of a neutral particle conserves the gauge symmetry of QED. In the context of the SM, as an effective theory, Majorana masses for neutrinos do not require new degrees of freedom. The gauge structure, however, demands the lowest dimensional operators for neutrino masses to be of dimension five [2]. These operators explicitly break the accidental lepton number global symmetry of the SM, which leads to lepton number violating (LNV) process, most notably the neutrino-less double beta decay ($0\nu\beta\beta$) of nuclei [3].

The Seesaw mechanism [4] represents a class of renormalizable models that UV complete the above setup of SM plus effective operators. In the simplest type-I seesaw model, gauge singlet right-handed (RH) neutrinos with large Majorana masses are integrated out at tree level to get the Weinberg operator $\lambda(L^T H)^2/M$. One can as well write down other dimension five operators, $\lambda(L^T\sigma^a L)(H^T\sigma^a H)/M$ or $\lambda(L^T\sigma^a H)(H^T\sigma^a L)/M$, which corresponds to integrating out at tree level a $SU(2)_L$ triplet scalar with hypercharge two or a triplet fermion without hypercharge, and are named type-II [5], III [6] seesaws, respectively. They can be embedded in well motived theories such as grand unified theories (GUT) [7]. Other incarnations of the seesaw mechanism include radiative seesaw which exists in the minimal supersymmetric model with R-parity violation [8, 9], and inverse seesaw which naturally arises from mirror symmetric models [10].

The phenomenological motivation for studying theories for the seesaw mechanism is to search for various LNV signals that are related to new physics responsible for seesaw. The seesaw scale M is of the most relevance to the energy frontier probes. If it lies within the accessibility of the Large Hadron Collider (LHC), we will have the opportunity to produce new particles which allow the direct probe of seesaw. On the other hand, neutrino mass requires $\lambda v^2/M \sim$ eV, where v is the electroweak scale. TeV seesaw scale M then means small coupling λ. Although this seems different from the original motivation where λ is of order one and M could be connected to the GUT scale, low scale seesaw is equally probable and even technically natural. At the intensity frontier, the seesaw sector make new contribution to $0\nu\beta\beta$, which can be described by the effective operator $(\bar{u}\Gamma_1 d)(\bar{u}\Gamma_2 d)(\bar{e}^c\Gamma_3 e)/M^5$ [11]. For TeV seesaw, the new contribution could give comparable signal strength to that from active neutrino masses.

The left-right symmetric model (LRSM) [12] has a profound theoretical motivation that connects the seesaw

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mechanism to the origin of parity symmetry violation. It extends the SM gauge group to a left-right symmetric one. The spontaneous gauge symmetry breaking from LRSM to SM violates parity in a maximal way, thus explains the large hierarchy between left (active) and RH neutrino mass scales. The minimal LRSM accommodates type-I and II seesaw mechanisms, and is arguably the simplest yet well motivated theory for neutrino masses. With the new right-handed charged and neutral currents gauge interactions, if LRSM is realized at TeV scale, it will offer a plethora of interesting phenomena (including LNV) at LHC [13, 14] and low energy experiments.

On the cosmology side, a question for the LRSM is the dark matter (DM) candidate. In the TeV-scale LRSM, the only viable candidate of DM is the lightest RH neutrino. Interestingly, in the minimal such model, there is no (approximate) Z_2 symmetry that could be defined for the (cosmological) stability of DM, which means the DM itself must be very light. Naively this is problematic because the fast gauge interactions will fully thermalize the RH neutrino DM and overproduce its relic density. The solution to this problem is to produce entropy in the universe after the freeze out of DM, and this could be done via the late decay of heavier RH neutrinos. A careful recent study in Ref. [15] reveals this dilution picture is compatible with a LRSM scale window of a few TeV. The key observation is the significant degrees of freedom change in the thermal universe during the QCD phase transition temperature. It coincides with the typical freeze out temperature of TeV scale gauge interactions, and offers the opportunity to develop a large number density hierarchy between the DM and the diluting particle. In the end, the DM mass is constrained to be around keV, which is a warm DM candidate.

This proceeding is organized as the following. In the next section, I first discuss the simplest type-II seesaw with double charged scalars, and the present constraints from the LHC and other related phenomena. In the second part of Sec. II, I discuss the type-I seesaw in the context of LRSM and the corresponding signatures. Sec. III is devoted to studying the cosmology of the LRSM, where I will focus on the DM candidate and show the unique (non-)thermal history for the correct DM relic density, and for the theory to be at the TeV scale.

TESTING SIMPLEST SEESAW THEORIES AT LHC

As part of the parameter space, TeV scale seesaw theories are interesting possibilities and are connected to a plethora of interesting signals at the LHC and low energy experiments. We discuss two such models: the type-II seesaw, and type-I seesaw when embedded in the context of LRSM. Through the gauge interactions, the new states for generating neutrino masses, if kinematically accessible, are friendly for direct production at the LHC.

Type-II Seesaw at LHC

In the type-II seesaw model, an $SU(2)_L$ triplet scalar with hypercharge $Y = 2$ is introduced,

$$\Delta = \begin{pmatrix} \Delta^+/\sqrt{2} & \Delta^{++} \\ \Delta^0 & -\Delta^+/\sqrt{2} \end{pmatrix} . \tag{1}$$

The most general Lagrangian for the scalar interactions is

$$\begin{aligned} \mathscr{L}_\Delta \;=\;& [(y_\Delta)_{ij} L_i^T C i\sigma_2 \Delta L_j + \text{h.c.}] + m_H^2 H^\dagger H - m_\Delta^2 \text{Tr}\Delta^\dagger\Delta + [\mu H^T i\sigma_2 \Delta^* H + \text{h.c.}] \\ & - \lambda_1 (H^\dagger H)^2 - \lambda_2 \text{Tr}(\Delta^\dagger\Delta)^2 - \lambda_3 \text{Tr}(\Delta^\dagger\Delta)^2 - \alpha H^\dagger H \,\text{Tr}\Delta^\dagger\Delta - \beta H^\dagger \Delta\Delta^\dagger H \,. \end{aligned} \tag{2}$$

The y_Δ and μ terms, in together with the usual charged lepton Yukawa coupling term, explicitly violate lepton number. At loop level, they may lead to processes such as $e_R^+ e_R^+ \to W^+ W^+$ or $0\nu\beta\beta$, although highly suppressed. However, without the Higgs vacuum expectation value (VEV), the gauge structure forbids neutrinos to pick up any Majorana mass.

After electroweak symmetry breaking, the Higgs VEV induces a triplet VEV, which breaks lepton number and generates neutrino masses at tree level,

$$(M_\nu)_{ij} = (y_\Delta)_{ij} v_\Delta \approx (y_\Delta)_{ij} \frac{\mu v^2}{m_{\Delta^0}^2} \,, \tag{3}$$

where in the last step Δ fields are integrated out at the neutrino mass scale, and $v_\Delta \equiv \langle \Delta^0 \rangle$. Electroweak precision measurement of the ρ parameter gives an upper bound on $v_\Delta \lesssim$ GeV.

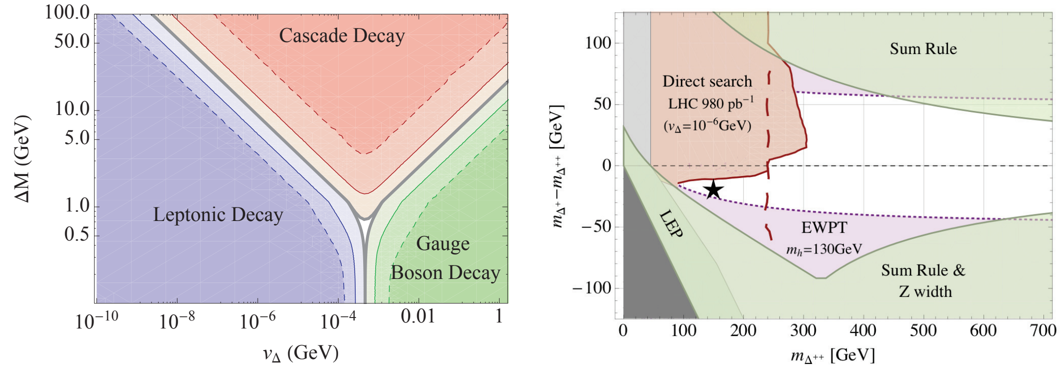

FIGURE 1. Left: Generic decay phase diagram for Δ decays in the type-II seesaw model, exemplified for case B defined in the text, with $m_{\Delta^{++}} = 150\,\text{GeV}$. Dashed, thin solid and thick solid contours correspond to 99, 90 and 50% of the branching ratios. Here $\Delta M = m_{\Delta^{++}} - m_{\Delta^+}$. Right: Summary of all the experimental and theoretical constraints in the $m_{\Delta^{++}}$–m_{Δ^+} parameter space, for degenerate light neutrino masses. The LHC 2σ exclusion is shown by the region to the left of the red solid curve, relative to $v_\Delta = 10^{-6}$ GeV. The analogous curve for $v_\Delta = 10^{-9}$ GeV is red dashed. The purple region is excluded by EWPT at 95% C.L. is shown for SM Higgs mass 130 GeV. The green region is excluded by the Z-width bound and the triplet mass sum rule.

We are interested in the case when the mass scales of Δ's are accessible at colliders. In this case, it is meaningful to examine the spectrum of Δ components. Because of the constraint $v_\Delta \ll v = 246\,\text{GeV}$, we will neglect the corrections at v_Δ order. The tree level masses are

$$m^2_{\Delta^{++}} = m^2_\Delta + \frac{1}{2}\alpha v^2, \;\; m^2_{\Delta^+} = m^2_\Delta + \frac{1}{2}\alpha v^2 + \frac{1}{4}\beta v^2, \;\; m^2_{\Delta^0} = m^2_\Delta + \frac{1}{2}\alpha v^2 + \frac{1}{2}\beta v^2 \,. \tag{4}$$

It is worth pointing out that they satisfy the approximate sum rule $m^2_{\Delta^+} - m^2_{\Delta^{++}} \simeq m^2_{\Delta^0} - m^2_{\Delta^+}$, i.e., the potential term β controls the universal splitting of mass squares.

The most striking signal at LHC is the decay of doubly-charged component Δ^{++} into two same-sign lepton final states [16, 17][18]. The direct searches using this channel and 7 and 8TeV LHC data can put very strong limit on the triplet mass scale. However, as pointed out in [16], this direct limit varies a lot depending on the parameters of this model (although simple). There are other possible decay channels that compete with the leptonic decay, as summarized below

$$\Delta^{++} \to \ell^+\ell^+, \, W^+W^+, \, \Delta^+W^+ \,. \tag{5}$$

The decay phase diagram is given in Fig. 1, which is controlled by only two parameters. If the v_Δ is large but β is small, Δ^{++} will decay into same-sign W-bosons, and the collider constraint based on the leptonic decay of W-boson is suppressed by branching ratio and becomes rather weak. It is even more so if the third decay channel dominates in which case Δ^+ will further cascade decay into $\Delta^0 W^+$, followed by $\Delta^0 \to \nu\nu$.

As an example, we show in the right panel of Fig. 1 the 7 TeV LHC limit on $m_{\Delta^{++}}$ in the general parameter space. We choose a small v_Δ so the competing channels are leptonic and cascade decays. When the mass difference vanishes $\beta = 0$, the leptonic decay has 100% branching ratio and the limit (red solid curve) agrees with that given by the CMS collaboration [19]. For $\beta > 0$, we have $m_{\Delta^0} > m_{\Delta^+} > m_{\Delta^{++}}$, and the limit using four lepton final states is enhanced because the productions of Δ^0 and Δ^+ will also contribute. On the other hand, when $\beta < 0$, the mass hierarchy is flipped, $m_{\Delta^0} < m_{\Delta^+} < m_{\Delta^{++}}$. Once produced, Δ^{++} will leak into lighter components and the leptonic decay of Δ^{++} is suppressed. In this case, the LHC limit decreases very fast with the mass splitting, to only 100 GeV or so when $m_{\Delta^{++}} - m_{\Delta^+} \gtrsim 20\,\text{GeV}$. With more up-to-date LHC data taken into account, the naive bound could improve, but the above parameter dependence will maintain.

One important implication of the limit on Δ^{++} mass scale has to do with the Higgs boson decay properties. If its mass is close to100 GeV, it can modify sizably the $h \to \gamma\gamma$ branching ratio, which also depends on the magnitude and sign of potential term α. For $\alpha < 0$, $h \to \gamma\gamma$ will be enhanced. This enhancement decouples quickly with the mass scale of Δ^{++}, and is negligible when $m_{\Delta^{++}} \gtrsim 300\,\text{GeV}$. Therefore, the precise direct limit plays a critical role here.

Type-I Seesaw at LHC via Left-Right Symmetry

In the case of Type-I seesaw, the RH neutrinos, N, are gauge singlets under the SM, which makes them hard to produce at colliders. The mixing angle between active and RH neutrinos is proportional to the neutrino Yukawa coupling, and typically suppressed by $\sqrt{m_\mu/m_N}$, which is 10^{-6} for weak scale RH neutrinos. It has been noticed [20], however, in the case of more than one generations, possible to have larger Yukawa elements, thus much larger mixing angle (~ 0.1) between light and heavy neutrinos, while maintaining the lightness of active neutrinos. It is thus possible to produce the RH neutrinos through weak interactions and probe its Majorana nature at colliders [21]. A critical comment on this scenario is that, this scenario is anyway fine-tuning the Yukawa couplings, and the flavor structure of $\ell_L - N$ weak interaction has little to do with the seesaw mechanism, i,e., the light neutrino mass and mixings.

The left-right symmetric model is based on the $SU(2)_L \times SU(2)_R \times U(1)_{B-L}$ gauge group (suppressing color), supplemented by the Parity symmetry between the left and the right sector. Quarks and leptons come in symmetric representations $Q_{L,R} = (u,d)^T_{L,R}$, $L_{L,R} = (\nu,\ell)^T_{L,R}$. Notice RH neutrinos must be introduced one for each generation for anomaly cancellation, and they are naturally charged under the new gauge symmetry. The generic form of the RH charged-current gauge interactions are

$$\mathscr{L}_{CC} = \frac{g}{\sqrt{2}} W_R^\mu \left[\begin{pmatrix} \overline{N_1} & \overline{N_2} & \overline{N_3} \end{pmatrix} \mathbf{V}_\ell^{R\dagger} \gamma_\mu P_R \begin{pmatrix} e \\ \mu \\ \tau \end{pmatrix} + \begin{pmatrix} \overline{u} & \overline{c} & \overline{t} \end{pmatrix} \mathbf{V}_q^R \gamma_\mu P_R \begin{pmatrix} d \\ s \\ b \end{pmatrix} \right] + \text{h.c.}, \tag{6}$$

This offers new prospects to produce the heavy RH neutrinos via the new gauge interactions. Their Majorana nature will be reflected at hadron colliders in the final states with same-sign dileptons, and without missing energy. Measuring the $\ell_R - N$ flavor structure in this case could be meaningful for understanding how seesaw works, for the structure of the LRSM connects left and right-handed neutrinos and their couplings. The Higgs sector of the minimal model consists of a bidoublet $\Phi = (2_L, 2_R, 0_{B-L})$ and two triplets, $\Delta_L = (3_L, 1_R, 2_{B-L})$ and $\Delta_R = (1_L, 3_R, 2_{B-L})$. The symmetry breaking in the model is characterized by the following VEV $\langle\Phi\rangle = \text{diag}(v_1, v_2)$, $\langle\Delta^0_{L,R}\rangle = v_{L,R}$, which have a hierarchical order $v_L^2 \ll v^2 = v_1^2 + v_2^2 \ll v_R^2$, with $v = 245$ GeV. The resulting masses of the heavy gauge bosons are $M_{W_R} = g\, v_R$ and $M_{Z_{LR}} \simeq \sqrt{3} M_{W_R}$. Notice also LRSM by construction includes the triplet Δ_L field, which is behind the type-II seesaw.

The presence of new RH charged current interactions can lead to new contribution to flavor-changing neutral current processes. In fact, the LRSM had been proposed [22] as a theory of CP violation in K-meson decays, without the need of third generation fermions. After the establishment of CKM theory of CP violation in the SM, the LRSM contribution to low-energy processes then translates into the lower bound on the scale of new gauge boson, M_{W_R}. The strongest bounds are from neutral K-meson mixing and CP violations, first pointed out in [23]. This bound depends crucially on the RH analog of the CKM matrix element between d and s quarks. In LRSM with general CP violation, the analytic result for the RH quark mixing matrix is found in [24] by expansions in m_b/m_t. This allows one to put solid bound on the RH scale, which is $M_{W_R} > 2.5$ TeV [24, 25] from K-meson mass difference and $M_{W_R} > 4$ TeV from CP violation observables [25]. A more recent analysis [26] confirmed this bound.

With the RH gauge interaction and RH Majorana neutrinos, there are also new sources contributions to the $0\nu\beta\beta$ process, in a left-right symmetric manner to the SM counterpart. At TeV scale, the RH contribution to $0\nu\beta\beta$ rate could dominate the SM one, and on the edge of current experimental discover potential.

On the LHC side, a pioneering study has showed the 14 TeV running will be able to access $M_{W_R} \sim 6-7$ TeV [27]. Compared with the low energy constraint, there is still a window remaining for probing the seesaw mechanism and the origin of Majorana masses via the left-right symmetry. This inspired us to perform a more detailed study [14] of the LHC limit and prospect using the early running data provided in [28]. The event topology we are interested in is

$$pp \to W_R^+ \to N_i \ell_j^+, \quad N_i \to \ell_k^\pm jj\,. \tag{7}$$

Because of the Majorana nature of RH neutrino N_i, there are equal probabilities for the two leptons ℓ_j, ℓ_k to be of same sign or opposite signs. Our results are shown in Fig. 2 in the joint parameter space of M_{W_R} and M_N. For RH neutrino mass in the range between 100 GeV and M_{W_R}, the above dilepton dijet channel is most effective, and it is clear that LHC limit already starts to exceed the constraint from low energy $0\nu\beta\beta$. For intermediate RH neutrino mass, it is more boosted and decay products will collimate each other. This reduces the efficiency of limit using the same channel. The future release of data on jets with electromagnetic activity is expect to put constraint on this case. For very light RH neutrinos, their lifetime is long enough to produce displaced vertex, or even be stable at collider time scale. In the latter case, the W_R mass limit coincides with that for a sequential W'. With the arrival of the latest 7 and 8 TeV LHC data, the above bounds in $2\ell 2j$ and $\ell + \not{E}_T$ channels have been significantly improved to around 2.9 TeV [29, 30].

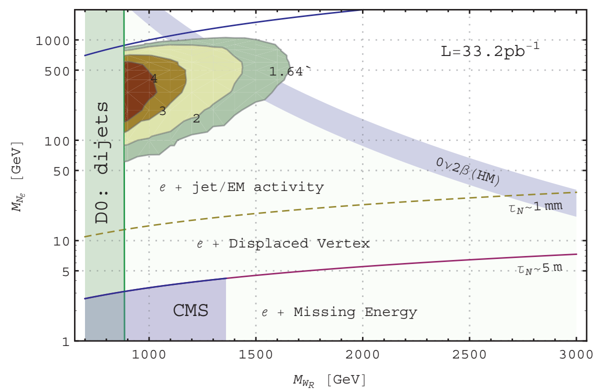

FIGURE 2. Combined limits in the M_{W_R}–m_N parameter space from the CMS data. The excluded vertical rectangular region on the left is the CMS result from $W_R \to tb$ decay. The excluded horizontal trapezoid below is the CMS missing energy result applicable for neglibly small m_N. For illustration, ignoring leptonic mixings, in the left plot we also depict a (grey) band where the LR contribution to $0\nu\beta\beta$ decay saturates the current limit.

LRSM AND DARK MATTER CANDIDATE

Based on the discussions in the previous section, it is clear that LRSM stands out as the theory behind the simplest seesaw model. In this section, we turn our focus to the dark side of the model, and examine whether it can also be a theory of dark matter. Of course, we are still interested in the low scale LRSM realized near TeV, so that all the phenomena discussed above are still within the reach of current and future experiments.

An interesting feature to notice first is that in the minimal LRSM, there is no exact or approximate enough Z_2 symmetry for the DM to be cosmologically stable. Therefore, the only possibility is the DM itself is light. There are two types of electric neutral particles introduced in the model: the RH neutrinos, and the "Higgs" boson $\mathrm{Re}\Delta_R^0$ responsible for the gauge symmetry breaking $U(1)_R \times U(1)_{B-L} \to U(1)_Y$. In analog with the SM, $\mathrm{Re}\Delta_R^0$ can decay into two photons via the RH gauge boson $W_R^\pm$. The decay rate may be sufficiently suppressed with $m_{\mathrm{Re}\Delta_R^0} \sim$keV and $M_{W_R} \gtrsim 10^{12}$TeV. This is against our interest of finding a TeV scale LRSM, and $m_{\mathrm{Re}\Delta_R^0} \ll M_{W_R}$ is also quite unnatural under radiative corrections, for a similar reason discussed in [31]. Therefore, the lightest RH neutrino N_1 is the only viable candidate.

The main obstacle we have to face when M_{W_R} lies in the TeV region, turns out to be the over-abundance of N_1, because the $SU(2)_R$ gauge interactions would fully thermalize it until the decoupling temperature T_f which is between a few hundred MeV to GeV. As discussed in Eq. (7), the decays $N_1 \to \ell jj$ or $N_1 \to \ell\pi$ are at collider time scale unless N_1 is much lighter than GeV. Therefore, we have $m_{N_1} \ll T_R$, and it is still a relativistic species during freeze out. The freeze out temperature is given by

$$T_f \simeq 400\,\mathrm{MeV}\left(\frac{g_*(T_f)}{70}\right)^{1/6}\left(\frac{M_{W_R}}{5\,\mathrm{TeV}}\right)^{4/3}. \tag{8}$$

Intuitively, one expects N_1 to have a number density similar to that of light neutrinos and overclose the universe, when combined with the most reliable cosmological lower limit on the DM mass, around a keV scale [32, 33, 34],

$$\Omega_{N_1} = \frac{Y_{N_1} m_{N_1} s}{\rho_c} \simeq 3.3 \times \left(\frac{m_{N_1}}{1\,\mathrm{keV}}\right)\left(\frac{70}{g_*(T_{f1})}\right). \tag{9}$$

This is to be contrasted with the latest Planck result [35] on the DM relic density $\Omega_{\mathrm{DM}} = 0.239 \pm 0.019$. Clearly the relic density is too large by more than one order of magnitude. But the bottomline here is the DM mass is pointed to keV range, which makes it a warm DM candidate.

The only way out of this impasse is to dilute the number density of N_1 by entropy production due to the late decay of some massive particle which dominates the universe [36]. Such a late decay should inject relativistic light SM particles that quickly equilibrate with the thermal plasma and "reheat" the photon temperature. In turn, it takes longer for the

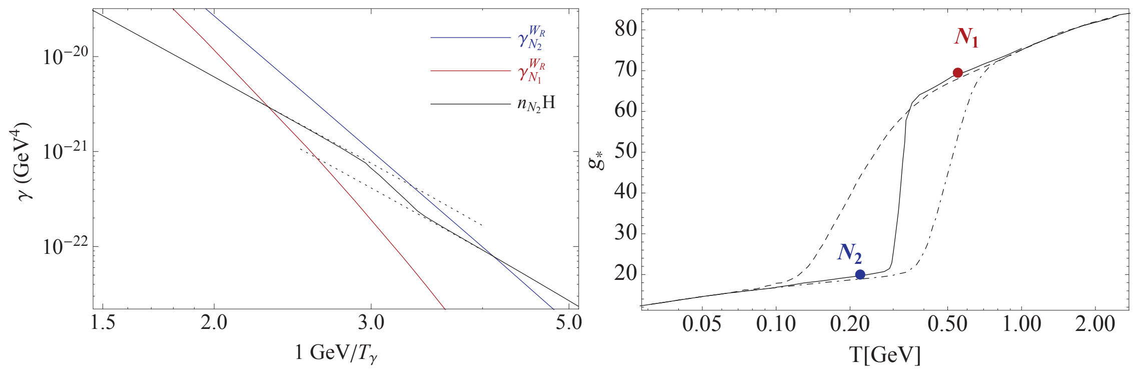

FIGURE 3. Left. Thermally averaged reaction rates for the processes that dominate the decoupling of N_1 (red) and N_2 (blue), for parameters $M_{W_R} = 5\,\text{TeV}$, $m_{N_2} = 0.25\,\text{GeV}$. Also shown is the Hubble expansion rate multiplied by the thermal number density of N_2 (black curve). **Right.** A sharp change in the evolution of g_{*S} around the QCD phase transition temperature. The dot-dashed (dashed) curve corresponds to $T_{\text{QCD}} = 400(150)\,\text{MeV}$ with a second-order QCD phase transition, while the solid line is an interpolation in between with $T_{\text{QCD}} = 350\,\text{MeV}$ and the transition close to first order. The red and blue points are the freeze-out temperatures of N_1 and N_2, respectively.

photons to cool down to present-day temperature and the number density of DM is effectively reduced. In order for the dilution to work, the temperature of N_1 should not increase, and it is therefore crucial that N_1 itself is not a decay product of the heavy decaying particle.

It has been shown [15, 37] in this model the heavier RH neutrinos $N_{2,3}$ are the only remaining viable candidates and they play the role of diluters in this scenario. Here, for the purpose of illustration of the generic picture and our idea, we proceed the discussion with a single diluter N_2. The generalization with more than one diluters will be shown with numerical results only. For calculation details we refer to [15]. As relativistic species, N_2 decouples with a yield

$$Y_{N_2} \equiv \frac{n_{N_2}}{s} \simeq \frac{135\,\zeta(3)}{4\pi^4\, g_*(T_{f_2})} \,. \tag{10}$$

If the lifetime of N_2 is long enough, as the temperature of the universe drops, after it becomes non-relativistic (and still abundant), sooner or later it will dominate the energy density of the universe. Then after its decay, the universe returns to the radiation dominant era. In the sudden decay approximation, energy conservation states the energy density of radiation after N_2 decay is equal to that of N_2 before the decay,

$$m_{N_2} n_{N_2} = m_{N_2} Y_{N_2} s_{before} = \rho_R(T_r) = \frac{3}{4} s_{after} T_r \,, \tag{11}$$

The reheating temperature T_r of the universe after the decay of N_2 is solely determined by its lifetime

$$T_r \simeq 0.78\, g_*(T_r)^{-1/4} \sqrt{\Gamma_{N_2} M_{\text{p}}} \simeq 1.22\,\text{MeV} \left(\frac{1\,\text{sec}}{\tau_{N_2}}\right)^{1/2} \,. \tag{12}$$

Therefore, the dilution factor, defined as the ratio of entropy before and after the sudden decay (no volume change), is given by

$$\mathscr{S} \equiv \frac{S_{\text{after}}}{S_{\text{before}}} \simeq \frac{s_{\text{after}}}{s_{\text{before}}} \simeq 1.8\,(g_*(T_r))^{1/4} \frac{Y_N\, m_N}{\sqrt{\Gamma_N M_{\text{p}}}} \,, \tag{13}$$

With this dilution, the relic density of N_1 will be relaxed to

$$\hat{\Omega}_{N_1} = \Omega_{N_1}/\mathscr{S} \simeq (0.239 + 0.019) \left(\frac{m_{N_1}}{1\,\text{keV}}\right) \left(\frac{1.91\,\text{GeV}}{m_{N_2}}\right) \left(\frac{1\,\text{sec}}{\tau_{N_2}}\right)^{1/2} \left(\frac{g_*(T_{f2})}{g_*(T_{f1})}\right) \,. \tag{14}$$

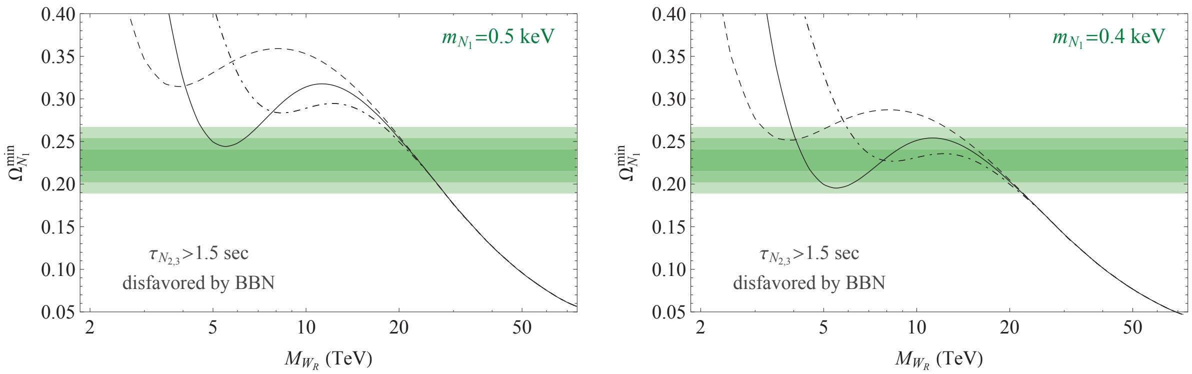

FIGURE 4. The minimum DM relic density that can be accommodated in the LR theory is plotted for two fixed DM masses as a function of M_{W_R}. Solid and dashed, dot-dashed lines correspond to $T_{\rm QCD} = 350$ MeV (nearly first order) and $T_{\rm QCD} = 150, 400$ MeV (second order), respectively and a fixed lifetime $\tau_{N_{2,3}} = 1.5$ sec. The green bands from dark to light correspond to DM relic abundance at 1, 2, 3 σ confidence level from WMAP fit. Keep in mind that the DM relic density is linearly proportional to m_{N_1}.

The next question is then whether the diluter lifetime can be made as long as 1 second. With a RH gauge interaction, a GeV scale RH neutrino decay rate is usually too short

$$\tau(N_2 \to \ell jj) \sim 10^{-8}\,{\rm sec} \left(\frac{2\,{\rm GeV}}{m_{N_2}}\right)^5 \left(\frac{M_{W_R}}{1\,{\rm TeV}}\right)^4 , \tag{15}$$

unless kinematically suppressed when $m_{N_2} \gtrsim m_\ell + m_\pi$. Take the $\ell = \mu$ case, the lifetime is long enough for the final state jets to hadronize into a pion, and the corresponding decay rate is

$$\tau(N_2 \to \mu\pi) \sim 1\ {\rm sec} \left(\frac{m_{N_2}}{250\,{\rm MeV}}\right)^{-3} \left(\frac{M_{W_R}}{5\,{\rm TeV}}\right)^4 \left(\frac{0.002}{f}\right) , \tag{16}$$

where f is a phase space factor, and can be suppressed when the final state phase space is kinematically squeezed. This offers a hint to make the lifetime long enough by tuning the diluter mass close to the threshold. Meanwhile, the flavor structure at the RH gauge vertex must be sufficiently diagonal so that there is only a single choice of lepton for N_2 to decay into.

Here we make a remark that the case $\ell = \tau$ is not a favored choice, although it seems the threshold $m_\tau + m_\pi$ fits very well with the desired value of diluter mass in Eq. (14). There are several reasons for this. First, such mass is too heavy compared to the typical freeze out temperature in Eq. (8). This will reduce the above dilution factor $\mathscr{S}$ by a Boltzmann suppression $e^{-(m_\tau+m_\pi)/T_{f_2}} \sim 10^{-2}$ and fail the dilution picture. Second, if N_2 couples mainly to τ, there is nothing to prevent the decay $N_2 \to \tau + N_1 + (e \text{ or } \mu)$, again through the RH gauge interaction. When the diluter can decay into DM (dilutee), the effect of dilution will be further suppressed. Based on these arguments, the desired flavor structure of RH neutrino gauge coupling (the RH PMNS matrix) emerges taking the form

$$\mathbf{V}_\ell^R \approx \begin{pmatrix} 0 & 0 & 1 \\ 0 & 1 & 0 \\ 1 & 0 & 0 \end{pmatrix}, \tag{17}$$

in together with the mass spectrum

$$m_{N_1} \sim \ {\rm keV}, \quad m_{N_2} \approx m_\pi + m_\mu, \quad m_{N_3} \approx m_\pi + m_e. \tag{18}$$

Because the diluter mass is now fixed to be lower than GeV, Eq. (14) is still in short by a factor of a few. As first pointed out in [15], it is possible to compensate for this factor using the hierarchy in the ratio $g_*(T_{f2})/g_*(T_{f1})$. The observation is that the freeze out temperature of RH neutrinos are of a few hundred MeV, which coincides with the QCD phase transition temperature, where the thermal dynamical nature of the universe, i.e., $g_*(T)$, changes drastically during a small temperature interval ΔT. Therefore, if the diluter freezes out at $T_{f_2} < T_{\rm QCD}$ while the dilutee freezes out at $T_{f_1} > T_{\rm QCD}$, we can get $g_*(T_{f2}) \ll g_*(T_{f1})$. The physical meaning of this is when the diluter freezes out at a bit lower temperature, it feels the reheating due to integrating out heavier hadron species, but the dilutee which freezes out earlier does not. In fact, the flavor structure is also designed for achieving such a situation. The point is τ lepton is becoming a heavy species at $T_{\rm QCD}$, and therefore the interaction maintaining the DM N_1 (which couples mainly to τ) in thermal equilibrium is slower compared to other flavors. This allows N_1 to freeze out relatively earlier. The interaction rates of N_i compared to the Hubble expansion rate is given in Fig. 3.

The interplay between TeV scale LRSM and the QCD phase transition in the early universe in this picture amounts to tuning the scale M_{W_R}, such that the relation $T_{f_2} < T_{\rm QCD} < T_{f_1}$ is fulfilled. The success of dilution then implies there is a TeV scale window of M_{W_R}, in which low scale LRSM is the theory of (warm) dark matter. This is shown in Fig. 4. In fact, this picture also survives from several important constraints, which we summarize in Table. 1. For more details, see [15], and [38].

TABLE 1. Various constraints on the masses and liftetimes of relevant states within the LRSM, coming from astrophysical, cosmological and terrestrial experiments, together with a sample point in the DM scenario.

Constraints	m_{N_1}	τ_N	M_{W_R}
Dwarf Galaxy	$\gtrsim 0.4-0.5\,\text{keV}$	—	—
Lyman-α	$\gtrsim 0.5-1\,\text{keV}$	—	—
BBN & CMB	—	$\lesssim 1.5\,\text{sec}$	—
$0\nu2\beta$	—	—	$\gtrsim 5-6\,\text{TeV}$
A sample point	$0.5\,\text{keV}$	$1.5\,\text{sec}$	$4-7\,\text{TeV}$

On the Fate of Baryon Number Asymmetry of the Universe

Before concluding, I briefly discuss the implication of the above pictures on the baryon asymmetry in the universe. A relistic history of the universe needs to explain not only the generation of DM but also the existence of the baryon asymmetry. Baryogenesis scenarios are the attempts from particle physics side to understand the generation of the baryon asymmetry. Although difficult to probe their existence directly, models for genesis at relatively low scale could lead to other related phenomena, therefore allow for interesting indirect experimental probes. Scenarios of this kind include electroweak baryogenesis (see [39] for a connection with the Higgs boson CP properties at LHC) or TeV leptogenesis. Unfortunately, it has been shown that TeV LRSM cannot be a theory of either [40]. This might make it less appealing. However, baryogenesis may as well happen from physics beyond LRSM. From an effective picture point of view, a less ambitious but more realistic question is whether the generated baryon asymmetry could survive, if LRSM interactions are operative, and what could be the qualified baryogenesis theories.

In the LRSM, it seems all the conserved quantum numbers related to baryon number would be broken, if the temperature of the universe had been as high as TeV scale. Above the electroweak scale, the $SU(2)_L$ weak sphaleron is in equilibrium and erases any B+L asymmetry. Meanwhile, if the RH gauge boson W_R exists in the plasma, the Majorana nature of RH neutrinos could washout all lepton number, via the process $e_R^+ e_R^+ \to W_R^+ W_R^+$. This process has the same topology as those for LHC production of RH neutrinos, and for $0\nu\beta\beta$. This is in spirit similar to the constraint on neutrino masses discussed in [41], but can be generalized to a broader one: If the universe has been hot enough, the survival of baryon asymmetry implies upper bound not only on the neutrino mass scale, but also on the rate of other lepton number violation processes that could be probed by present experiments. Therefore, I conclude here that if the LRSM is realized at TeV scale and LNV signals are to be seen in the near future, a viable cosmological history for baryogenesis must happen at a much lower scale.

SUMMARY AND OUTLOOK

To summarize, I have discussed the theories of seesaw mechanism for neutrino mass, and their phenomenology at LHC and low energy experiments. The LHC (non-)observation of the LNV processes already updated the bound on these theories. I discussed the blind spots and prospects of LHC searches in specific models. I argue the left-right symmetry model to be the theory of neutrino mass, which has its own motivations, and can accommodate the simplest version of seesaw. I also show that LRSM in a TeV scale window can be a viable theory of warm DM candidate, with the light RH neutrino being the DM. This will keep being probed by the current and future experiments at different frontiers of particle physics.

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