

## Is leptogenesis falsifiable at LHC?

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JHEP01(2009)051

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## Is leptogenesis falsifiable at LHC?

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ABSTRACT: It is well known that the leptogenesis mechanism offers an attractive possibility to explain the baryon asymmetry of the universe. Its particular robustness however comes with one major difficulty: it will be very hard if not impossible to *test* experimentally in a foreseeable future, as most of the mechanics *typically* takes place at high energy or *results* from suppressed interactions, without unavoidable low-energy implications. An alternate approach is taken by asking: can it be at least falsified? We show that *possible* discoveries at current and future colliders, most notably that of right-handed gauge interactions, would indeed forbid at least the "canonical" leptogenesis mechanisms, namely those based on right-handed neutrino decay. General lower bounds for successful leptogenesis on the mass of the right-handed gauge boson  $W_R$  are given. Other possibilities to falsify leptogenesis, including from the observation of a  $Z'$ , are also considered.

KEYWORDS: Baryogenesis, Beyond Standard Model, Neutrino Physics, Cosmology of Theories beyond the SM.

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## 1. Introduction

The recent evidence for neutrino masses has brought forward leptogenesis [1] as a very attractive mechanism to explain the baryon asymmetry of the universe. Along this mechanism, the baryon asymmetry of the universe is explained by the same interactions as the ones which can explain the neutrino masses. In the most straightforward seesaw model, which assumes right-handed neutrinos in addition to the standard model particles, both neutrino masses and leptogenesis originate from the Yukawa interactions and lepton number violating Majorana masses of the right-handed neutrinos

$$\mathcal{L} \ni -\bar{L} \tilde{H} Y_\nu^\dagger N - \frac{1}{2} \bar{N} m_N N^c + \text{h.c.} \quad (1.1)$$

where  $L$  stands for the lepton weak doublets and  $\tilde{H}$  is related to the standard Brout-Englert-Higgs (hereafter simply Higgs) doublet  $H \equiv (H^+, H^0)$  by  $\tilde{H} = i\tau_2 H^*$ .

However, testing this mechanism will be a very difficult task for several reasons. If the right-handed neutrinos have a hierarchical mass spectrum, due to neutrino mass constraints, leptogenesis through  $N$  decay can lead to the observed amount of baryon asymmetry e.g. only if it involves right-handed neutrinos with masses above  $\sim 10^8$  GeV [2, 3]. As a result they cannot be produced at colliders. Moreover there are many more parameters in the Yukawa coupling matrices which can play an important role for leptogenesis, than there are (not too suppressed) low energy observables which could constrain these parameters.<sup>1</sup>

If the right-handed neutrinos have instead a quasi-degenerate spectrum (for at least 2 of them), leptogenesis can be efficient at lower scales [5] but generically in this case the neutrino mass constraints require suppressed values of Yukawa couplings, which hampers their production at colliders.

For leptogenesis to be both efficient and tested at low energy, not only is a quasi-degeneracy between 2 right-handed neutrinos required, but also a special flavour structure which allows for larger Yukawa couplings while preserving the light neutrino mass constraints,<sup>2</sup> and/or a right-handed neutrino production mechanisms other than through the Yukawas and associated neutrino mixings.

In this paper we consider the problem of testing leptogenesis mechanisms the other way around. While they cannot confirm leptogenesis, could low energy observations at least exclude it? We propose one particularly clear possibility, namely the observation of a right-handed charged gauge boson  $W_R$ . It is known that for high mass right-handed neutrinos and  $W_R$ , around  $10^{10}$  GeV or higher, the  $W_R$  can have suppression effects on leptogenesis through dilution and scattering, but, in the specific case of reheating after inflation, they can also boost the  $N$  abundances [10–12] and hence relax the constraints on Yukawa couplings. Not surprisingly, with a low scale  $W_R$  the suppression effects are dramatically enhanced. Actually, see section 2, they turn out to be so strongly enhanced that, even with a maximal CP asymmetry of order unity, leptogenesis cannot be a sufficient cause of the matter excess anymore.

Right-handed gauge interactions lead in particular to much larger suppression effects at low scale than left-handed interactions do in other contexts (i.e. than in leptogenesis from scalar [13, 14] or fermion [15] triplet decays, whose efficiency have been calculated in refs. [14, 15]). This is due to the fact that at the difference of triplets, a single  $N$  can interact through  $W_R$  exchange with fermions which are all in thermal equilibrium, which

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<sup>1</sup>A possible exception to that arises in the supersymmetric case from the effects of Yukawa couplings on the running of the slepton masses [4]. This nevertheless assumes that universality of lepton soft mass terms must be present (an assumption which requires to be tested) and, for any real test of leptogenesis, would require to observe a long series of rare leptonic decays not necessarily expected to be all close to the present corresponding experimental bounds.

<sup>2</sup>This case can be realized if the Yukawa induced dimension 6 operator coefficients are unsuppressed (decoupling from the suppressed neutrino mass dimension 5 ones). This does not necessarily require cancellations of the various entries. It requires that some of entries are smaller than others, as in the inverse seesaw, see e.g. [7, 6, 8]. But it e.g. leads only to lepton conserving channels with rather large background at LHC [9].

induces more efficient, and hence dangerous, scatterings and decays. In particular, some of the scatterings involving the  $W_R$  turn out to induce a very large suppression due to the fact that they do not decouple through a Boltzmann suppression. The production of  $N$ 's through a light  $W_R$ , often presented as the easiest way to produce  $N$ 's, is therefore incompatible with successful leptogenesis, and even enhanced  $N$  production from reheating cannot compensate for the large suppression. The lower bounds on the mass of the  $W_R$ , required for successful leptogenesis, are given in section 3.

The possible discovery of a low-energy  $W_R$  has recently been the object of several analysis by LHC collaborations [16–18]. It should be feasible up to  $m_{W_R} \sim 3\text{-}5\text{ TeV}$  (see more details, and additional possible searches, in section 7).

The observation of a  $W_R$  is not the only possibility to exclude canonical neutrino decay leptogenesis from current energy data. We give a list of other possibilities in section 5, considering in particular the implications of the observation of a  $Z'$  at LHC. The case of other leptogenesis seesaw models with not only or without right-handed neutrinos is briefly considered in section 6.

## 2. Leptogenesis in presence of a low scale $W_R$

As well known the net rate of baryon asymmetry is given in any leptogenesis model by 3 ingredients, the CP asymmetry of the decaying particle,  $\varepsilon_N$  for a right-handed neutrino, the Boltzmann equations which determine the efficiency  $\eta$  and the  $L$  to  $B$  sphaleron conversion rate, which we denote by  $r_{\mathcal{L}\rightarrow\mathcal{B}}$ . Let us first discuss and present our results for the case where the lepton asymmetry is created from the decay of a single right-handed neutrino,  $N$ .<sup>3</sup> Later on we will discuss the generalization to more right-handed neutrinos. In this case, from these 3 ingredients the net baryon asymmetry produced by the  $N$  decays is:

$$Y_B = Y_{\mathcal{L}} r_{\mathcal{L}\rightarrow\mathcal{B}} = \varepsilon_N \eta Y_N^{eq}(T \gg m_N) r_{\mathcal{L}\rightarrow\mathcal{B}}. \quad (2.1)$$

with  $Y_i \equiv n_i/s$ ,  $Y_B \equiv Y_B - Y_{\bar{B}}$ ,  $Y_{\mathcal{L}} \equiv Y_L - Y_{\bar{L}}$ ,  $n_i$  the comoving number density of the species "i", "eq" referring to the equilibrium number density, and  $s$  the comoving entropy density. For a particle previously in thermal equilibrium, the efficiency is unity by definition in absence of any washout effect from inverse decays or scatterings. If all lepton asymmetry has been produced before the sphaleron decoupling at the electroweak phase transition and if the sphalerons have had the time to thermalize completely the  $L$  abundance, the conversion ratio between lepton and baryon number is given by [19]

$$r_{\mathcal{L}\rightarrow\mathcal{B}} = -\frac{8n_f + 4n_H}{22n_f + 13n_H} = -\frac{28}{79}, \quad (2.2)$$

where the last equality refers to the SM value, with  $n_f$  the number of fermion families and  $n_H$  the number of Higgs doublets.

In the right-handed neutrino decay leptogenesis model without any  $W_R$ , the CP-asymmetry is defined by

$$\varepsilon_N \equiv \frac{\Gamma(N \rightarrow LH) - \Gamma(N \rightarrow \bar{L}H^*)}{\Gamma(N \rightarrow LH) + \Gamma(N \rightarrow \bar{L}H^*)}. \quad (2.3)$$

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<sup>3</sup>We will not consider finite temperature effects which are not expected to change our conclusions.

while the evolution of the comoving abundances is given as a function of  $z \equiv m_N/T$  by the Boltzmann equations:

$$zH(z)s Y'_N = - \left( \frac{Y_N}{Y_N^{\text{eq}}} - 1 \right) \left( \gamma_N^{(l)} + 2\gamma_{Hs} + 4\gamma_{Ht} \right) \quad (2.4)$$

$$zH(z)s Y'_{\mathcal{L}} = \gamma_N^{(l)} \left[ \varepsilon_N \left( \frac{Y_N}{Y_N^{\text{eq}}} - 1 \right) - \frac{Y_{\mathcal{L}}}{2Y_L^{\text{eq}}} \right] - 2 \frac{Y_{\mathcal{L}}}{Y_L^{\text{eq}}} \left( \gamma_{Ns}^{\text{sub}} + \gamma_{Nt} + \gamma_{Ht} + \gamma_{Hs} \frac{Y_N}{Y_N^{\text{eq}}} \right) \quad (2.5)$$

where ' denotes the derivative with respect to  $z$ . The thermally averaged reaction rate

$$\gamma_N^{(l)} = n_N^{\text{eq}}(z) \frac{K_1(z)}{K_2(z)} \Gamma_N^{(l)}, \quad (2.6)$$

parametrizes the effects of Yukawa induced decays and inverse decays with  $\Gamma_N^{(l)} = \Gamma(N \rightarrow LH) + \Gamma(N \rightarrow \bar{L}H^*) = \frac{1}{8\pi} |Y_\nu|^2 m_N$ , and  $K_{1,2}$  Bessel functions. The other  $\gamma$ 's take into account the effects of the various scatterings through a  $H$  or a  $N$  in the  $s$  or  $t$  channels. They are related to the corresponding cross sections in the following way

$$\gamma(ab \leftrightarrow 12) = \iint d\bar{p}_a d\bar{p}_b f_a^{\text{eq}} f_b^{\text{eq}} \iint d\bar{p}_1 d\bar{p}_2 (2\pi)^4 \delta^4(p_a + p_b - p_1 - p_2) |\mathcal{M}|^2 \quad (2.7)$$

$$= \frac{T}{64 \pi^4} \int_{s_{\text{min}}}^{\infty} ds \sqrt{s} \hat{\sigma}(s) K_1 \left( \frac{\sqrt{s}}{T} \right) \quad (2.8)$$

with  $\hat{\sigma} = 2s^{-1} \lambda^2 [s, m_a^2, m_b^2] \sigma(s)$  the reduced cross section,  $\lambda[a, b, c] \equiv \sqrt{(a-b-c)^2 - 4bc}$  and  $s_{\text{min}} = \max[(m_a + m_b)^2, (m_1 + m_2)^2]$ . The analytic expression of the reduced cross sections can be found in refs. [20, 21].<sup>4</sup>  $\gamma_{Ns}^{\text{sub}} = \gamma_{Ns} - \gamma_N^{(l)}/4$  in eq. (2.5) refers to the subtracted scattering through a  $N$  in the  $s$  channel (i.e. taking out the contribution of the on-shell propagator in order to avoid double counting with the inverse decay contribution [21]).

The above, now traditional approach assumes that  $N$  are introduced in an isolated way in the model. In many unifying groups (left-right symmetric [22], Pati-Salam [23], SO(10) [24] or larger) the presence of the  $N$  can be nicely justified as it is precisely the ingredient required to unify all fermions. These groups however do not introduce the  $N$  in such an isolated way and moreover link the  $N$  and  $W_R$  masses to the same SU(2)<sub>R</sub> breaking scale  $v_R$ .<sup>5</sup> It is thus a (generally unwarranted) assumption to neglect the effect of SU(2)<sub>R</sub> gauge bosons. If  $m_{W_R}$  is smaller than  $\sim 10^{13}$  GeV, these effects must be explicitly incorporated for any  $N$  whose mass is not several orders of magnitude below the one of the  $W_R$  [11].

The key interactions of the  $W_R$  [22, 23] are the

$$\mathcal{L} \ni \frac{g}{\sqrt{2}} W_R^\mu (\bar{u}_R \gamma_\mu d_R + \bar{N} \gamma_\mu l_R) \quad (2.9)$$

<sup>4</sup>Note that for simplicity we have neglected the subdominant effects of scatterings of the type  $N + L \leftrightarrow H + (\gamma, Z, W_L)$  [21]. We also neglect as in ref. [21] the effects of Yukawa coupling induced  $NN \leftrightarrow LL, HH$  processes which have little effects too.

<sup>5</sup>More complicated breaking mechanisms could add extra contributions to the gauge boson masses: all mass contributions to  $N$  will also contribute to  $W_R$ , but the opposite is not necessarily true.

gauge ones. ( $N$  and the right-handed charged leptons ( $l_R = e_R, \mu_R, \tau_R$ ), and  $u_R$  and  $d_R$ , are members of a same  $SU(2)_R$  doublet).

Their effects for leptogenesis can be incorporated by modifying the Boltzmann equations in the following way:

$$zH(z)s Y'_N = - \left( \frac{Y_N}{Y_N^{\text{eq}}} - 1 \right) \left( \gamma_N^{(l)} + \gamma_N^{(W_R)} + 2\gamma_{Hs} + 4\gamma_{Ht} + 2\gamma_{Nu} + 2\gamma_{Nd} + 2\gamma_{Ne} \right) - \left( \left( \frac{Y_N}{Y_N^{\text{eq}}} \right)^2 - 1 \right) \gamma_{NN} \quad (2.10)$$

$$zH(z)s Y'_L = \gamma_N^{(l)} \varepsilon_N \left( \frac{Y_N}{Y_N^{\text{eq}}} - 1 \right) - \left( \gamma_N^{(l)} + \gamma_N^{(W_R)} \right) \frac{Y_L}{2Y_L^{\text{eq}}} - \frac{Y_L}{Y_L^{\text{eq}}} \left( 2\gamma_{Ns}^{\text{sub}} + 2\gamma_{Nt} + 2\gamma_{Ht} + 2\gamma_{Hs} \frac{Y_N}{Y_N^{\text{eq}}} + \gamma_{Nu} + \gamma_{Nd} + \gamma_{Ne} \frac{Y_N}{Y_N^{\text{eq}}} \right) \quad (2.11)$$

with the CP asymmetry unchanged, as given by eq. (2.3). In these Boltzmann equations there are essentially 2 types of effects induced by the  $W_R$ , both suppressing the produced lepton asymmetry: from the presence of alternate decay channels for the heavy neutrinos,  $\gamma_N^{(W_R)}$ , and from scatterings,  $\gamma_{Nu,d,e}$ , see below.

### 2.1 Decay effect: dilution and wash-out

It is useful to distinguish 2 cases depending on the mass hierarchy between  $N$  and  $W_R$ .

- a) Case  $m_{W_R} > m_N$ : in this case the decay of  $N$  to leptons or antileptons plus Higgs particles remains the only possible 2 body decay channels but a series of three body decay channels with a *virtual*  $W_R$  is now possible:  $N \rightarrow l_R q_R \bar{q}'_R$  or  $N \rightarrow \bar{l}_R \bar{q}_R q'_R$  with  $l = e, \mu, \tau, q = u, c, t, q' = d, s, b$ . We obtain:

$$\Gamma(N \rightarrow l_R q_R \bar{q}'_R) = \frac{3 g_R^4}{2^9 \pi^3 m_N^3} \int_0^{m_N^2} dm_{12}^2 \frac{(m_N^6 - 3m_N^2 m_{12}^4 + 2m_{12}^6)}{(m_{W_R}^2 - m_{12}^2)^2 + m_{W_R}^2 \Gamma_{W_R}^2(m_{12}^2)} \quad (2.12)$$

Given the potentially large value of the gauge to Yukawa couplings ratio, the three body decays can compete with the Yukawa two body decay. Since the gauge interactions do not provide any CP-violation and are flavor blind, it can be shown that they do not provide any new relevant source of CP-asymmetry. But still the gauge interaction-induced 3 body decays appear in both Boltzmann equations, eqs. (2.10)–(2.11), with

$$\gamma_N^{(W_R)} = n_N^{\text{eq}}(z) \frac{K_1(z)}{K_2(z)} \Gamma_N^{(W_R)}. \quad (2.13)$$

where  $\Gamma_N^{(W_R)}$  is the total three body decay width.

Unlike in leptogenesis without  $W_R$ , not all decays participate in the creation of the asymmetry but only a fraction  $\Gamma_N^{(l)}/\Gamma_{N_{\text{Tot}}}$  does. This shows up in the Boltzmann equations through the fact that eq. (2.10) involves  $\Gamma_{N_{\text{Tot}}} = \Gamma_N^{(l)} + \Gamma_N^{(W_R)}$  while the

CP-asymmetry in eq. (2.11) is multiplied only by  $\Gamma_N^{(l)}$ .<sup>6</sup> This dilution effect leads automatically to an upper bound on the efficiency. The bound  $\eta < 1$ , which applies in standard leptogenesis for thermal  $N$ 's becomes:

$$\eta < \Gamma_N^{(l)}/\Gamma_{N_{\text{Tot}}} \quad (2.14)$$

As a numerical example, for  $m_N \sim 1$  TeV, with Yukawa couplings of order  $10^{-6}$ , so that  $m_\nu \sim Y_\nu^2 v^2/m_N \sim 10^{-1}$  eV, and with  $m_{W_R} \sim 3(4)$  TeV we obtain the large suppression factor  $\Gamma_N^{(l)}/\Gamma_{N_{\text{Tot}}} = 7 \cdot 10^{-7}(2 \cdot 10^{-6})$ , consistent with leptogenesis only if the CP-asymmetry is of order unity, which requires maximal enhancement of the asymmetry (i.e. right handed neutrino mass splittings of order of their decay widths).

In addition to this dilution effect, the three body decay  $\gamma_N^{(W_R)}$  reaction density also induces a  $L$  asymmetry washout effect from inverse decays (proportional to  $Y_{\mathcal{L}}$  in eq. (2.11)) which can also be large.

- b) Case  $m_{W_R} < m_N$ : in this case<sup>7</sup> the direct 2 body decays  $N \rightarrow W_R l_R$  are allowed which leads to an even larger dilution and washout effect for low  $m_N$ . For example with  $m_N \simeq 1$  TeV,  $Y_\nu \simeq 10^{-6}$  and  $m_{W_R} \simeq 800$  GeV, we get  $\Gamma_N^{(l)}/\Gamma_{N_{\text{Tot}}} = 4 \cdot 10^{-9}$ , which means that the dilution effect makes leptogenesis basically hopeless at this scale, even with the maximum value  $\varepsilon_N = 1$ . In the following we will consider only the case where  $m_{W_R} \gtrsim m_N$  (this corresponds to the situation where a discovery of the  $W_R$  and  $N$  at LHC would occur through same sign dilepton channel [16, 17, 25], see section 6).

## 2.2 Gauge scattering effect

Right-handed gauge interactions induce a long series of scatterings, given in figure 1. To explain their effects let us first consider scatterings which do not involve any external  $W_R$ , figure 1a. The density reaction rates  $\gamma_{Nu}$ ,  $\gamma_{Nd}$ ,  $\gamma_{Ne}$ ,  $\gamma_{NN}$  can be computed from the following reduced cross sections:

$$\hat{\sigma}(Ne_R \rightarrow \bar{u}_R d_R) = \frac{9g_R^4}{8\pi s[(s - m_{W_R}^2)^2 + m_{W_R}^2 \Gamma_{W_R}^2]} \left( \frac{m_N^6}{6} - \frac{m_N^2 s^2}{2} + \frac{s^3}{3} \right) \quad (2.15)$$

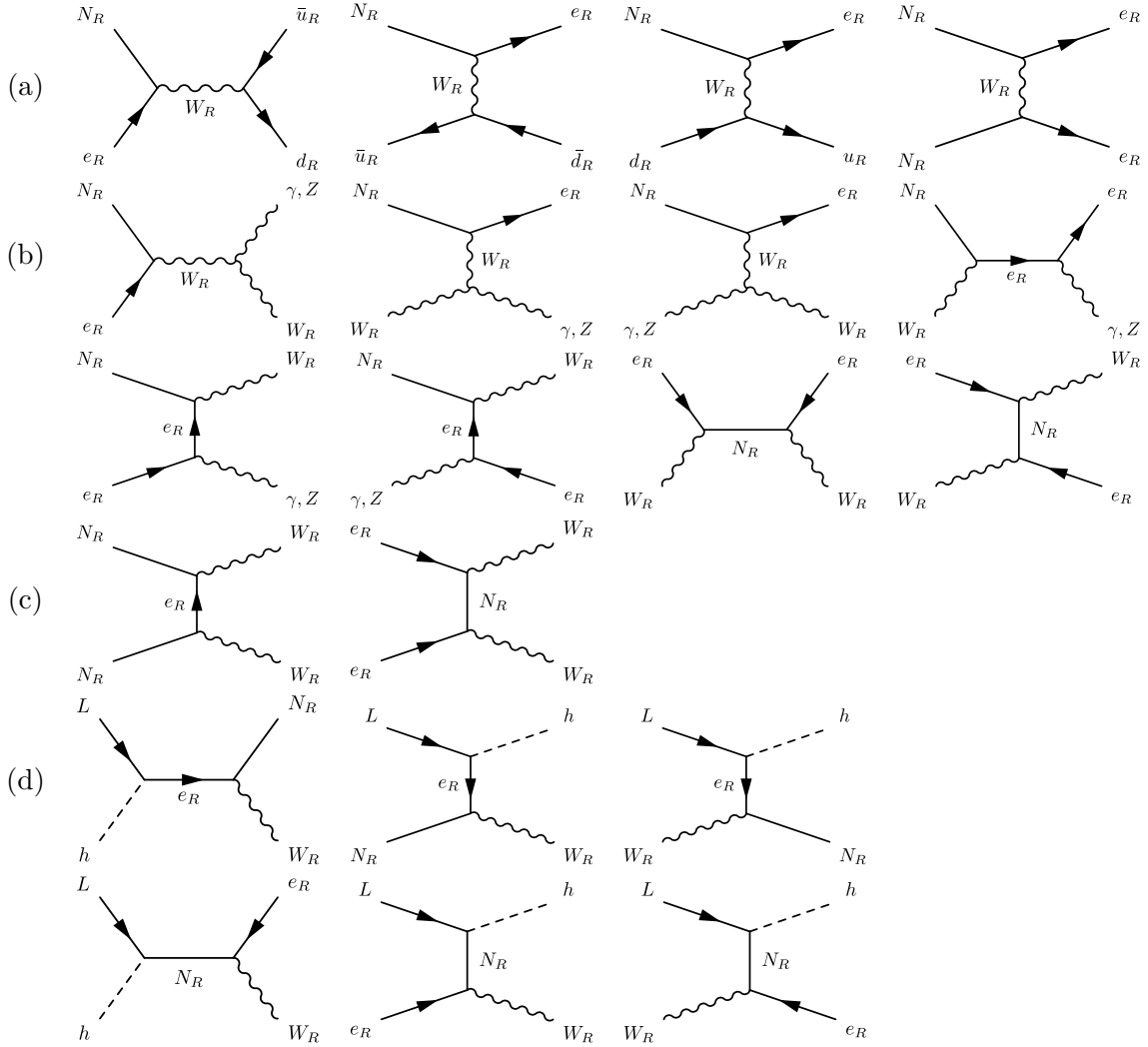
$$\hat{\sigma}(N\bar{u}_R \rightarrow e_R \bar{d}_R) = \frac{9g_R^4}{8\pi s} \int_{m_N^2 - s}^0 dt \frac{(s+t)(s+t - m_N^2)}{(t - m_{W_R}^2)^2} \quad (2.16)$$

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<sup>6</sup>In eq. (2.11), we made the choice to keep eq. (2.3) as definition for the CP-asymmetry. In its denominator, it involves only the Yukawa driven decay rather than the total decay width,  $\Gamma_{N_{\text{Tot}}}$ . Therefore this CP asymmetry doesn't correspond anymore, as in standard leptogenesis, to the averaged  $\Delta L$  which is created each time a  $N$  decays. However this definition is convenient for several reasons. It makes explicit the fact that the gauge decay does not induce any lepton asymmetry. Moreover in this way, all (competing) suppression effects, including the dilution one, are put together in the efficiency, not in the CP-asymmetry. It also allows to take the simple upper bound  $\varepsilon < 1$  for any numerical calculations.

<sup>7</sup>A  $N$  much heavier than  $W_R$  is in general not expected in the left-right symmetric model or extensions given the fact that, as said above, both  $W_R$  and  $N_R$  have a mass proportional to the  $SU(2)_R$  breaking scale  $v_R$ , and given the fact that  $m_{W_R} \sim gv_R$  with  $g$  the ordinary gauge coupling which is of order unity.





**Figure 1:** Scatterings involving the  $W_R$ .

$$\hat{\sigma}(Nd_R \rightarrow e_R u_R) = \frac{9g_R^4}{8\pi} \frac{(m_N^2 - s)^2}{m_{W_R}^2 (s + m_{W_R}^2 - m_N^2)} \quad (2.17)$$

$$\hat{\sigma}(NN \rightarrow e_R \bar{e}_R) = \frac{g_R^4}{8\pi s} \int_{t_0}^{t_1} dt \left( \frac{(s + t - m_N^2)^2}{(t - m_{W_R}^2)^2} + \frac{(m_N^2 - t)^2}{(2m_N^2 - s - t - m_{W_R}^2)^2} - \frac{m_N^2 s}{(t - m_{W_R}^2)(2m_N^2 - s - t - m_{W_R}^2)} \right) \quad (2.18)$$

Among these scatterings the three first ones involving only one external  $N$  have a peculiar property. Unlike in ordinary pair annihilation or in coannihilation with a heavier particle, their decoupling in the  $Y_N$  Boltzmann equation does not proceed with a Boltzmann suppression of their rate. The decoupling condition is:

$$\frac{\gamma_A}{n_N^{eq} H} \lesssim 1 \quad (2.19)$$

with  $H$  the Hubble constant and  $\gamma_A = \gamma_{Nu} + \gamma_{Nd} + \gamma_{Ne}$ . For  $T$  well below  $m_N$  the reaction density, eq. (2.8), is Boltzmann suppressed (i.e. in  $e^{-m_N/T}$ ) but so is also  $n_N^{eq}$  in the denominator. Therefore, decoupling comes at low temperature only from the approximately linear in  $T$  behaviour of  $\frac{\gamma_A}{n_N^{eq}H}$  for small  $T$ . This can be understood from the fact that what sets the thermal equilibrium of  $Y_N$  is the number of interactions per  $N$ , not the number of interactions irrespective of the number of  $N$ . In other words these processes are important because the abundance of the other particles involved is large with respect to the  $N$  density.

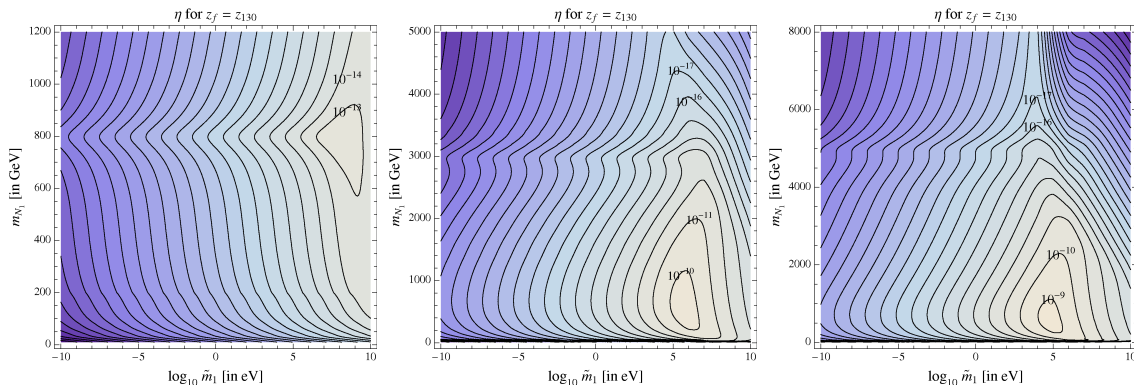
It is useful to compare this behaviour with the one of ordinary left-handed gauge scatterings which have been considered for leptogenesis from the decay of a scalar triplet [14] or of a fermion triplet [15]. In these models these scatterings necessarily involve two external heavy-states (i.e. annihilation or creation of a pair of scalar triplets or a pair of fermion triplet respectively) and therefore are doubly Boltzmann suppressed (which leads to an exponential Boltzmann type decoupling:  $\frac{\gamma}{n_T^{eq}H} \sim e^{-m_T/T}$ ).

The right-handed gauge interaction induced scatterings remain therefore in thermal equilibrium down to temperatures much lower than the left-handed gauge triplet interactions for equal decaying state and gauge boson masses. Their decoupling also doesn't occur so sharply (compare for example  $\gamma_A$  with  $\gamma_{NN}$  in figure 3 below or with the left-handed gauge scattering rates of figure 3 of ref. [14] or of figure 6 of ref. [15]).

For  $m_{W_R}$  and  $m_N$  of order TeV, one observes from a numerical analysis that the decoupling temperature which follows from eq. (2.19) is  $\sim 15$  orders of magnitude below these masses. At this temperature the number of  $N$  remaining is hugely Boltzmann suppressed, so that no sizeable asymmetry can be created. However, due to the fact that their decoupling is not sharp, these scatterings still allow the creation of a highly suppressed but non-vanishing lepton asymmetry at temperature well above this value (see numerical results below). In all cases the later the  $N$  decays with respect to  $m_{W_R}$ , the less the gauge scatterings will be in thermal equilibrium at the time of the decays, and the smaller will be the suppression effect from them.

Note also that unlike the left-handed gauge interactions, the suppressions from the scatterings of eqs. (2.15)–(2.17) also operate in the  $Y_{\mathcal{L}}$  Boltzmann equation, eq. (2.11). This can lead to several orders of magnitude further suppression (see below). The decoupling of these scatterings in the  $Y_{\mathcal{L}}$  Boltzmann equation results from a Boltzmann suppression when  $\gamma_A/(n_i^{eq}H) \lesssim 1$ . In ref. [11] these effects of gauge scatterings (as well as of three body inverse decays) in the  $Y_{\mathcal{L}}$  Boltzmann equation have been omitted. In the region of parameters considered in this reference, these effects are nevertheless moderate, see below.

Beside the gauge scattering of figure 1a there are also scatterings with one external  $W_R$  changing the number of  $N$  and/or violating lepton number, figure 1b. Since a substantial asymmetry can be created only at temperature as low as possible, well below  $m_{W_R}$  for  $m_{W_R} \gtrsim m_N$ , all these scatterings are suppressed with respect to the ones with no external  $W_R$ , eqs. (2.15)–(2.17). The relative suppression effect is  $e^{-m_{W_R}/m_N}$ . Similarly the scatterings with two external  $W_R$ , figure 1c are further suppressed. Finally the scatterings of figure 1d are suppressed by powers of the Yukawa couplings. As a result we will neglect all



**Figure 2:** For values of the right-handed gauge boson mass which could be probed at LHC,  $m_{W_R} = 0.8, 3, 5$  TeV, iso-efficiency curves as a function of  $\tilde{m}$  and  $m_N$ . As expected the efficiency decreases from right to left panel, and is always too suppressed to obtain successful baryogenesis.

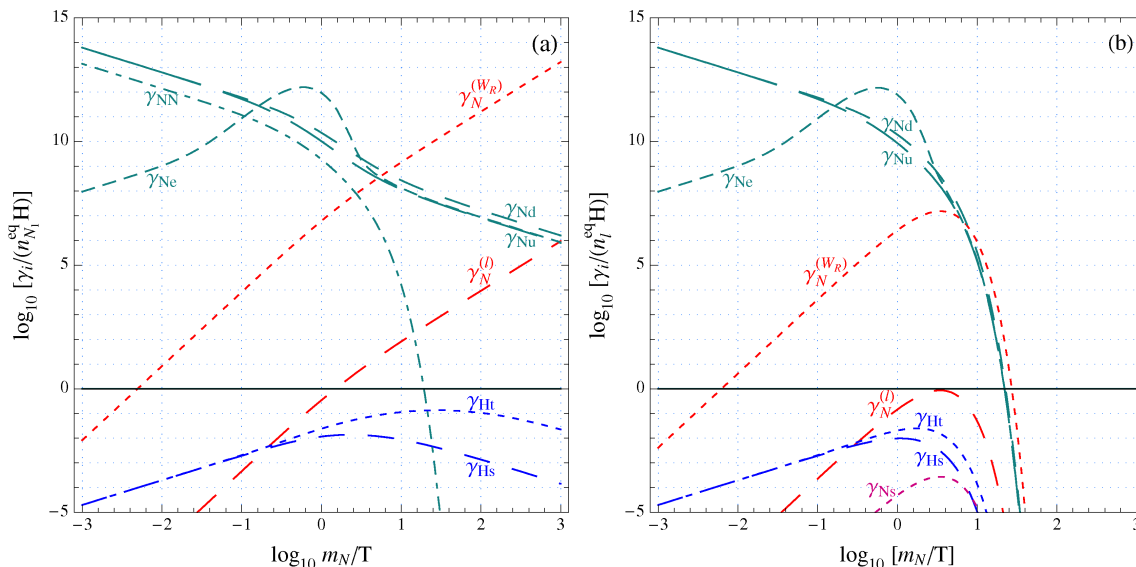
the scatterings of figure 1b–1d and keep only the ones of figure 1a.<sup>8</sup>

### 2.3 Efficiency results

All in all the efficiency we obtain numerically is given in figure 2, as a function of  $m_N$  and  $\tilde{m} = v^2 Y_\nu^\dagger Y_\nu / m_N = \Gamma_N^{(l)} 8\pi v^2 / m_N^2$  for various values of  $m_{W_R} = 800$  GeV, 3 TeV, 5 TeV with  $v = 174$  GeV.  $m_{W_R} = 800$  GeV corresponds essentially to the lower experimental limit [26], while  $m_{W_R} = 3$  TeV corresponds essentially to the value LHC could reasonably reach [18]. Motivated by the analysis of ref. [27], these figures are based on the approximation that all  $L$  asymmetry produced above  $T \sim 130$  GeV (for  $m_h \sim 120$  GeV) has been converted to a  $B$  asymmetry (with conversion factor as given in eq. (2.2)), but none of it afterwards. In all cases we get an efficiency factor far below  $\sim 7 \cdot 10^{-8}$  which is the minimum value necessary to get the observed baryon asymmetry  $Y_B = (6 - 9) \cdot 10^{-11}$  (with maximal CP-asymmetry).

To understand these results it is useful to discuss the effect of the various terms step by step. For this, we take as example the set of parameters:  $m_N = 500$  GeV,  $m_{W_R} = 3$  TeV,  $\tilde{m} = 10^{-3}$  eV. figure 3 provides the various reaction densities divided by  $n_N^{eq} H$  and  $n_l^{eq} H$ , as relevant for discussing thermal equilibrium in the  $Y_N$  and  $Y_L$  Boltzmann equation respectively. Figure 4 gives the  $Y_N$  and  $Y_L$  abundances as a function of  $z$ . As well known, omitting all  $W_R$  interactions, figure 4a, there is no large efficiency suppression for  $\tilde{m} = 10^{-3}$  eV, we get  $\eta \simeq 0.5$ , i.e.  $Y_B = 6.2 \cdot 10^{-4}$  (with  $\varepsilon_N = 1$ ). Adding to this case only the effect of the 3 body decay in the  $Y_N$  Boltzmann equation, figure 4b, leads to the dilution effect explained above:  $\eta \simeq \gamma_N^{(l)} / \gamma_N^{(W_R)} \simeq 2.8 \cdot 10^{-8}$ , i.e.  $Y_B \simeq 3.6 \cdot 10^{-11}$ . Adding the gauge scattering terms in the  $Y_N$  Boltzmann equation leads to a even more suppressed result for any  $z < 6.5$  because in this range  $\gamma_A > \gamma_N^{(W_R)}$ . Given the fact that the sphaleron decoupling temperature corresponds to  $z \simeq 4$  we do get an extra suppression:  $\eta \simeq 1.5 \cdot 10^{-10}$ , i.e.  $Y_B \simeq 1.8 \cdot 10^{-13}$ , figure 4c. The efficiency is roughly given by the

<sup>8</sup>These scatterings can only further suppress leptogenesis, which as we will see is anyway already far too suppressed to be successful.

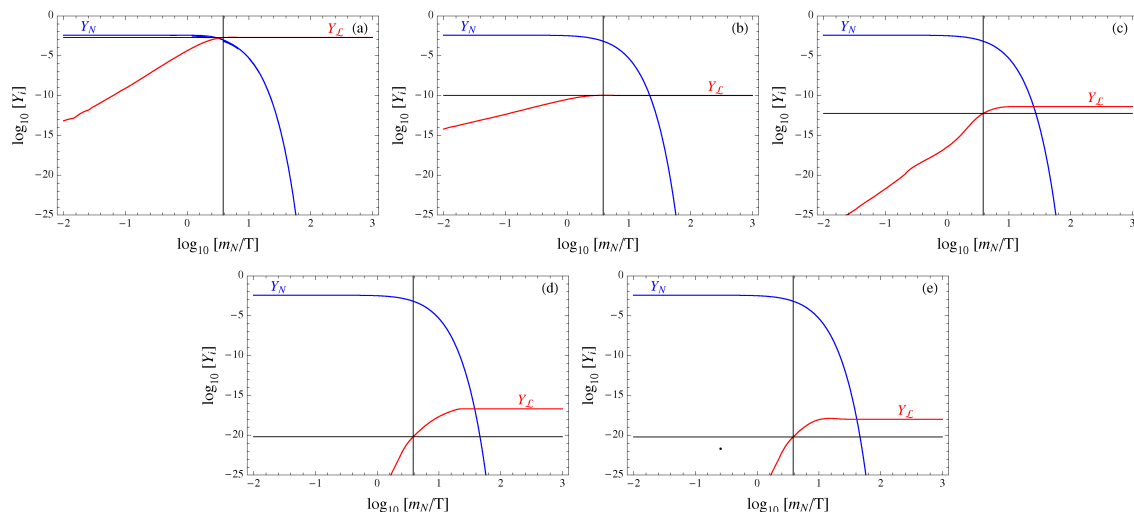


**Figure 3:** Evolution of the reaction densities (a)  $\gamma/(n_N^{eq}H)$  and (b)  $\gamma/(n_l^{eq}H)$  with  $z$  for  $m_N = 500$  GeV,  $m_{W_R} = 3$  TeV and  $\tilde{m} = 10^{-3}$  eV.

value of  $\gamma_A/\gamma_N^{(l)}$  a bit before sphaleron decoupling. Note that the result is sensitive to the sphaleron decoupling temperature. For smaller decoupling temperatures where  $\gamma_A$  is smaller the efficiency would have been larger and would have lead to about the same result as in figure 4b. Adding furthermore the  $\Delta L = 1$  gauge scattering effects in the  $Y_C$  Boltzmann equation, figure 4d, leads to further suppression because for  $T > 130$  GeV, these scatterings turn out to be fast enough to put leptons close to chemical equilibrium, i.e.  $\gamma_A/n_l^{eq}H > 1$ , see figure 3b. We get:  $\eta \simeq 1.6 \cdot 10^{-18}$ , i.e.  $Y_B \simeq 2.1 \cdot 10^{-21}$ . Finally adding the 3 body decay effect to the  $Y_C$  Boltzmann equation doesn't lead to further sizable suppression at  $T = 130$  GeV because above this temperature  $\gamma_A > \gamma_N^{(W_R)}$ . Only between  $z \simeq 6.5$  (when  $\gamma_N^{(W_R)}$  becomes larger than  $\gamma_A$ ) and  $z = 30$  (when  $\gamma_N^{(W_R)}/n_l^{eq}H$  becomes smaller than 1) it could have had an effect, compare figure 4d and figure 4e. Altogether at  $T = 130$  GeV we get  $\eta \simeq 1.6 \cdot 10^{-18}$  as given in figure 2.

Note that for  $m_{W_R} = 3$  TeV, the values  $m_N \simeq 500$  GeV and  $\tilde{m} \simeq 10^5$  eV appear to be the ones which maximize the efficiency. Larger values of  $m_N$  lead to more suppression from the  $W_R$ . Smaller values lead to a creation of the asymmetry occurring too late to be converted by the sphalerons. The important effect of sphaleron decoupling for low  $N$  mass can be seen by comparing figure 2b with figure 5 where no sphaleron decoupling temperature cut has been applied. Similarly smaller values of  $\tilde{m}$  leads to more suppressed efficiency from larger  $\gamma_A/\gamma_N^{(l)}$  and  $\gamma_N^{(W_R)}/\gamma_N^{(l)}$  ratios in the  $Y_N$  Boltzmann equation. Large values of  $\tilde{m}$  lead though to very large suppression from Yukawa driven inverse decays and  $\Delta L = 2$  scatterings. Those effects start to dominate over the  $W_R$  effects for  $\tilde{m} \simeq 10^5$  eV, which explains why in figure 2a maximum is got around this value of  $\tilde{m}$ :  $\eta \simeq 10^{-10}$ .

Note also that, for  $m_N \sim m_{W_R}$ , in figure 2, there is a local enhancement of the efficiency because, as  $m_N$  approaches  $m_{W_R}$  from below, the  $\gamma_A$  rate becomes more and more



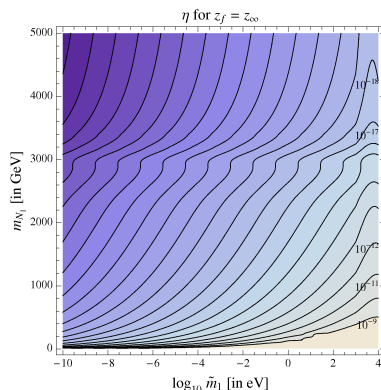
**Figure 4:** Evolution of  $Y_N$  and  $Y_{\mathcal{L}}$  abundances as a function of  $z = m_N/T$  for  $m_N = 500$  GeV,  $m_{W_R} = 3$  TeV and  $\tilde{m} = 10^{-3}$  eV, including various effects in the Boltzmann equations as explained in the text. The straight lines indicate the value of  $z$  and  $Y_{\mathcal{L}}$  at sphaleron decoupling.

insensitive to the  $W_R$  resonance. However as  $m_N$  gets larger than  $m_{W_R}$  the  $N \rightarrow W_R l_R$  decay opens up and the efficiency gets again suppressed.

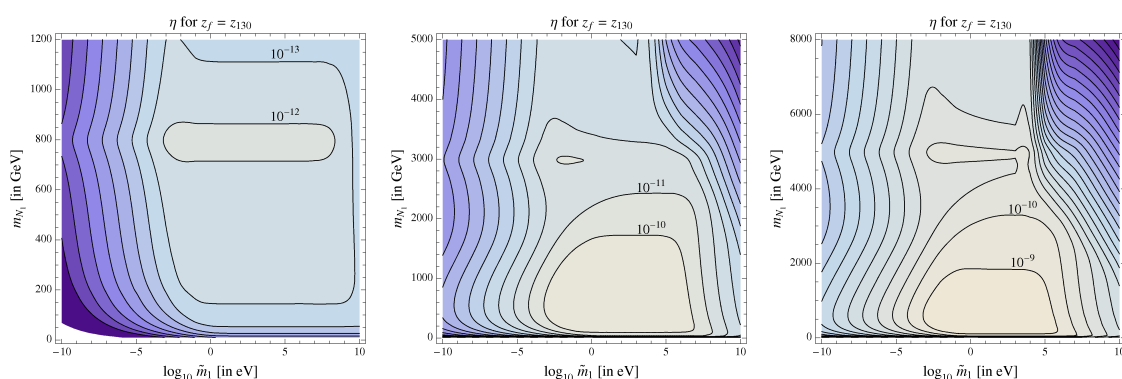
One additional question one must ask is whether our results depend on the fact that we considered only the evolution of the total lepton number asymmetry. The results can indeed largely depend on the flavour structures of the Yukawa couplings as well as on the flavour of the  $SU(2)_R$  light partner of the  $N$ , but not enough to allow successful leptogenesis. For example even if  $N$  could create an asymmetry only in flavours orthogonal to the flavour of its  $SU(2)_R$  partner, leptogenesis still wouldn't work. In this case the asymmetry produced wouldn't be washed-out by any  $W_R$  interaction appearing in the  $Y_{\mathcal{L}}$  Boltzmann equation, but still the  $W_R$  thermalization effects in the  $Y_N$  Boltzmann equation would be fully effective since they do not depend on flavour.<sup>9</sup> We have checked over the full  $\tilde{m}$  and  $m_N$  parameter space that even in this extreme case we would get a far too suppressed efficiency to have successful leptogenesis. Our results for this case are given in figure 6, see also the example of figure 4.

One more question to ask is whether the results obtained above could sizeably depend on the initial distribution of  $N$  before they decay. The answer is simply no, due to the fact that, starting from any number of  $N$  at temperature above  $m_N$  (from no  $N$  to only  $N$  in the universe) the  $W_R$  interactions very quickly put the  $N$ 's in deep thermal equilibrium.

Note finally that since we neglected the scatterings of figure 1b and figure 1c, strictly speaking our result is valid only for  $m_N < m_{W_R}$ . But this is where the maximum efficiency is obtained and elsewhere these scatterings can only suppress even more leptogenesis.



**Figure 5:** Efficiencies without sphaleron decoupling for  $m_{W_R} = 3$  TeV. (For values of  $\tilde{m}$  beyond  $10^5$  eV, Yukawa driven  $\Delta L = 2$  scatterings are so fast that the efficiency collapses.)



**Figure 6:** Iso-efficiency curves for  $m_{W_R} = 0.8, 3, 5$  TeV as a function of  $\tilde{m}$  and  $m_N$  when gauge interactions are only present in the  $Y_N$  Boltzmann equation.

### 3. Bounds on $m_{W_R}$ and $m_N$

In the previous section we have seen that for  $m_{W_R}$  reachable at LHC, successful leptogenesis from  $N$  decays is not possible. Larger values of  $m_{W_R}$  lead however to better efficiencies. It is useful to determine what are the bounds on  $m_{W_R}$  for a given value of  $m_N$  and vice versa. These can be obtained from figure 7a which for fixed values of  $m_{W_R}$  gives the allowed range of  $m_N$  and  $\tilde{m}$  taking the maximum value  $\varepsilon_N = 1$ . One observes that the absolute lower bound on  $m_{W_R}$  is 18 TeV. It is obtained for  $m_N = 500$  GeV and  $\tilde{m} = 3 \cdot 10^2$  eV. This value of  $\tilde{m}$  requires large cancellations between large Yukawa couplings in the neutrino masses. More usual values lead to a more severe bounds, we get

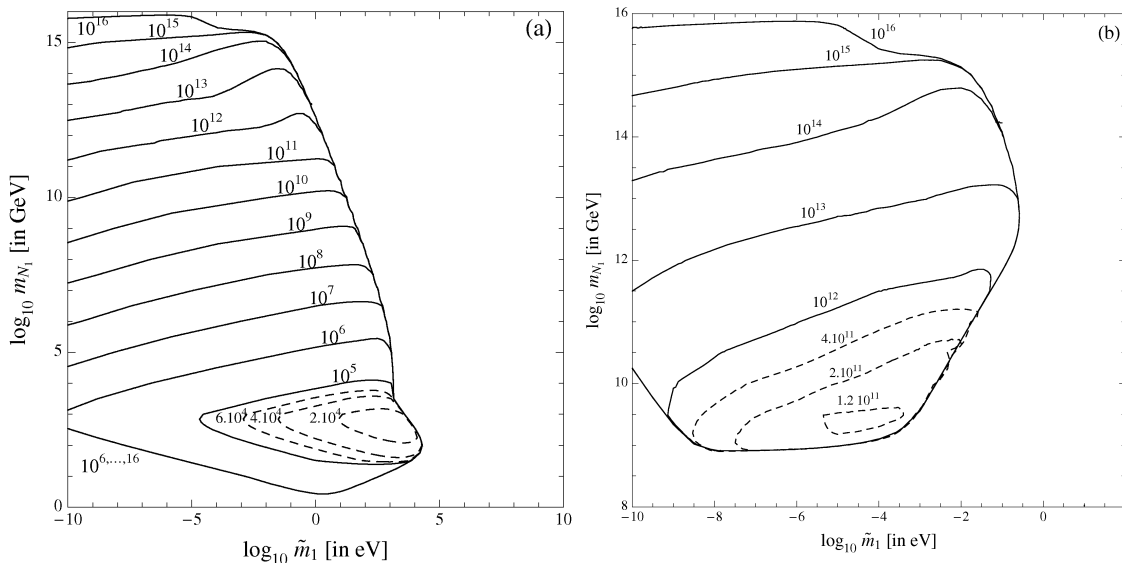
$$m_{W_R} > 110, 60, 35 \text{ TeV} \quad \text{for} \quad \tilde{m} = 10^{-5, -3, -1} \text{ eV} \quad (3.1)$$

we get the bound

$$m_N > 2.6 \text{ GeV} \quad (3.2)$$

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<sup>9</sup>We neglect effects of charged leptons Yukawa couplings which are much less important.



**Figure 7:** For various values of  $m_{W_R}$  (in GeV), the inner part of each curve gives the values of  $\tilde{m}$  and  $m_N$  which can lead to successful leptogenesis (i.e.  $Y_{\mathcal{B}} = 9 \cdot 10^{-11}$ ). Left (right) panel is obtained for  $\varepsilon_N = 1$  ( $(3/16\pi) m_N \sqrt{\Delta m_{atm}^2}/v^2$ ). The dependance in  $m_{W_R}$  of the lower bound on  $m_N$  is totally negligible, except for  $m_{W_R} < 10^6$  (left panel) and  $m_{W_R} < 2 \cdot 10^{11}$  (right panel).

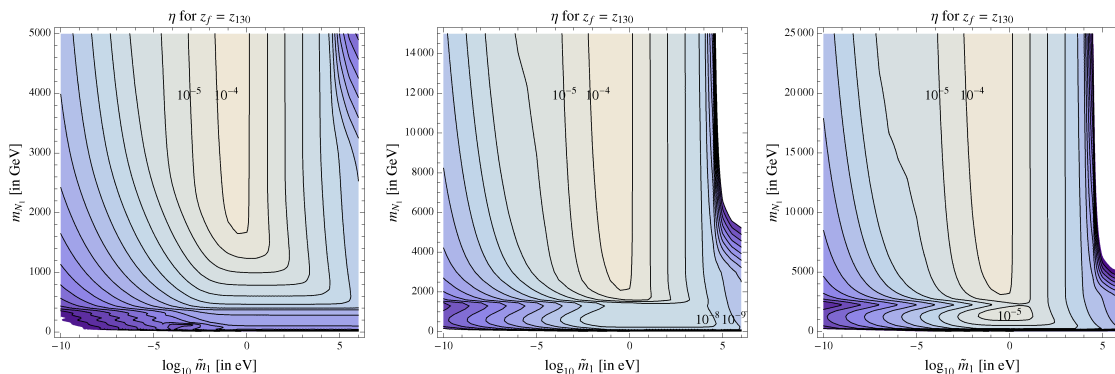
which holds even for the case where  $W_R$  effects are negligible. This gives an absolute lower bound on  $m_N$  which is another tantalizing target for excluding leptogenesis.

For completeness we also give in figure 7b the results we obtain taking the lower bound  $\varepsilon_N < (3/16\pi) m_N \sqrt{\Delta m_{atm}^2}/v^2$  [3] which holds for a hierarchical spectrum of right-handed neutrinos. We obtain the absolute bound  $m_{W_R} > 10^{11}$  GeV which requires  $m_N = 2.6 \cdot 10^9$  GeV and  $\tilde{m} = 5 \cdot 10^{-5}$  eV. We also get

$$m_{W_R} > 1.1 \cdot 10^{11}, 1.3 \cdot 10^{11}, 1.1 \cdot 10^{12} \text{ GeV} \quad \text{for} \quad \tilde{m} = 10^{-5, -3, -1} \text{ eV}. \quad (3.3)$$

The flavour dependance of the results of this section is relatively moderate. For the extreme case above where all  $W_R$  have been omitted in the  $Y_{\mathcal{L}}$  Boltzmann equation, instead of equation eq. (3.1), we get  $m_{W_R} > 39, 13, 8.8$  TeV, while the absolute lower bound on  $m_{W_R}$  becomes 8.7 TeV which we obtained for  $\tilde{m} = 10^1$  eV. The bounds of eq. (3.3) in this case are relaxed by less than 10 percent, while the lower bounds on  $m_N$ , as well as the upper bounds on  $\tilde{m}$ , are negligibly affected in figures 7a and 7b. As for the upper bounds on  $m_N$  in these figures, they are relaxed by up to one order of magnitude. The results of figure 7b agree with the one of [11] for what can be compared, modulo these flavour effects, since the  $W_R$  effects are neglected in the  $Y_{\mathcal{L}}$  Boltzmann equation in this reference.

Note that we do not expect that the results of figure 7 could be largely affected by the (neglected) scatterings of figure 1b-c, because all bounds in these figures are obtained with  $m_N \lesssim m_{W_R}$  (except in corners of parameters space for large  $m_{W_R}$  and large  $\tilde{m}$  where it is not excluded that these scatterings could reduce the bounds on  $m_N$  by up to a few times).



**Figure 8:** Iso-efficiency curves for  $m_{Z'} = 0.8, 3, 5$  TeV as a function of  $\tilde{m}$  and  $m_N$ .

#### 4. Generalization to several right-handed neutrinos

The results obtained above are strictly valid only if the lepton asymmetry is produced by a single right-handed neutrino, the effects of the other heavy states being present only in the CP asymmetry  $\varepsilon_N$  and in the  $\Delta L = 2$  washout.<sup>10</sup> Consequently these results assume that the heavier states do not create their own asymmetry and do not induce any washout besides this  $\Delta L = 2$  one. However, we are not aware of any model where  $\varepsilon_N$  can be obtained as large as unity, the upper bound we considered above, and where the above assumption can be fully justified. For instance, as said above, one possibility to have large CP asymmetries at low scale is through quasi-degeneracy of at least 2 right-handed neutrinos leading to a resonant enhancement of the self-energy diagram. In this case to a very good approximation both right-handed neutrinos have equal CP-asymmetries and equal masses, which means that both  $N_{1,2}$  must be considered in the Boltzmann equations. In the appendix A we show that this does not change though our conclusions. The point is that the asymmetry produced by two neutrinos is bounded by the sum of both asymmetries we get in the single  $N$  case with  $\tilde{m} = \tilde{m}_1$  and with  $\tilde{m} = \tilde{m}_2$  (with  $\tilde{m}_i$  referring to the value of  $\tilde{m}$  of  $N_i$ ), eq. (A.8). From the results of figures 2 and 6 this shows that the lepton asymmetry produced will be always too small to produce enough asymmetry if  $m_{W_R}$  is as low as in these figures, as relevant for the LHC. Furthermore from this inequality, if both  $\tilde{m}_i$  lie outside the range of values allowed by figure 7a, a large enough baryon asymmetry cannot be produced. Moreover it can be checked numerically that this figure remains also valid to a good approximation for the case  $\tilde{m} = \tilde{m}_1 = \tilde{m}_2$ . It is in this sense that this figure has to be interpreted for the several  $N$  case.

#### 5. Other possible suppression effects

##### 5.1 Effects of a $Z'$ associated to a U(1) symmetry

A  $Z'$  associated to an extra low energy U(1) could be discovered at LHC up to  $\sim 3$ -

<sup>10</sup>In  $\gamma_{N_s}^{sub}$  and  $\gamma_{N_t}$  above we took into account the contributions from  $N_{2,3}$  proportional to the neutrino masses, as given in eqs. (92, 93) of ref. [21] with  $\xi = \sqrt{\Delta m_{atm}^2}/\tilde{m}$ , because these contributions are relevant anyway (even for hierarchical  $N$ 's) for very large  $m_N$  and/or very large  $\tilde{m}$ .



5 TeV [25, 17]. If it couples to  $N$  through the  $Z'^{\mu}(\bar{N}\gamma_{\mu}N)$  interaction it has effect on the efficiency through the  $Y_N$  Boltzmann equation. Since this interaction involves 2  $N$  it doesn't induce any relevant 2 or 3 body decays which could cause dilution, and the associated scatterings decouple through a Boltzmann suppression. As a result the suppression effect is not as large as with a  $W_R$ . For example considering a  $U(1)_{Y'}$  as it has been considered in [28], see also [29], including all associated scatterings (i.e. the effect of  $NN \leftrightarrow f\bar{f}$ ,  $HH$  scatterings), the efficiency we obtain for  $M_{Z'} = 0.8, 3, 5$  TeV is given in figure 8. It shows that the discovery of a  $Z'$  would not necessarily rule out leptogenesis depending on the values of  $\tilde{m}$ , but would require very large values of  $\varepsilon_N$ .

### 5.2 Effects of a $Z'$ associated to a $SU(2)_R$ symmetry

The neutral gauge boson associated to  $SU(2)_R$  symmetry could also be discovered at LHC up to  $\sim 3$ -5 TeV [25, 17]. Since it is in the same multiplet as the  $W_R$ , its effect should be included in the analysis above together with the effects of the  $W_R$ . As it also couples only to 2  $N$ , the suppression effects due to this neutral gauge boson will nevertheless be negligible with respect to the ones of the  $W_R$  when the asymmetry is created: the  $N$  will have an interaction involving a  $W_R$  before having one involving the  $Z'$  (as long as  $m_{Z'} \simeq m_{W_R}$  as expected in the left-right symmetric models).

### 5.3 Effects of a right-handed triplet

The consequences of the discovery of one or several components of a right-handed scalar triplet  $\Delta_R = (\delta_R^{++}, \delta_R^+, \delta_R^0)$  could be dramatic for leptogenesis in some cases.

The easiest state to discover at LHC is the doubly charged one,  $\delta_R^{++}$ , due to suppressed background in the same sign dilepton channel [31]. As this state couples only to 2 right-handed charged leptons [22], and doesn't couple directly to the  $N$ , it has no sizable effect on the  $Y_N$  Boltzmann equation but can have an effect on the second one through L-violating  $l_R l_R H H$  interactions mediated by the  $\delta_R^{++}$ . This effect can be large if the couplings involved are of order  $\sim 10^{-4}$  or larger depending on the masses. The presence of the  $\delta_R^{++}$  would be however indicative of the existence of other triplet members.

A  $\delta_R^+$  (e.g. more difficult to see at LHC because it doesn't produce same sign dilepton channels in as direct a way as the  $\delta_R^{++}$ ), can couple to a  $N$  and a  $l_R$  as the  $W_R$ . It can therefore induce dilution effect from the  $N \rightarrow \delta_R^+ l_R$  decay if kinematically allowed, or from  $N \rightarrow l_R H^+ H^0$  decays otherwise (i.e. through a  $\delta_R^+ H^- H^0$  coupling with  $H$  any lighter scalar particle, e.g. from the bidoublet in LR models [22]). Similarly it induces dangerous scatterings similar to the one of figure 1a, replacing the  $W_R$  by a  $\delta_R$  and the quark pair by a  $H^+ H^0$  pair. For couplings in these processes as large as the  $W_R$  gauge couplings, the suppression of the efficiency is expected to be similar to the one caused by the  $W_R$  in section 2, which would rule out leptogenesis. For smaller couplings however the suppression decreases quickly. In the later case leptogenesis can be successfully produced from  $N \rightarrow \delta_R^+ l_R$  decays if kinematically allowed [32].<sup>11</sup>

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<sup>11</sup>The observation of a  $W_R$  would rule out this leptogenesis mechanism in the same way as in section 2.

Finally the  $\delta_R^0$  couples to  $2 N$  and therefore is expected to have effects roughly similar to the ones of a  $Z'$ , if the Yukawa couplings are as large as the gauge couplings, less otherwise.

#### 5.4 Effects of a neutral or charged $SU(2)_L$ scalar singlet

In large varieties of models, e.g. non left-right, a  $SU(2)_L$  scalar singlet can couple to  $2 N$  if it is neutral or to a  $N$  and a  $e_R$  if its electromagnetic charge is unity. These states, if they also couple to right-handed quarks, can be dangerous for leptogenesis in a similar way as the above  $\delta_R^0$  and  $\delta_R^+$  states respectively.

### 6. Suppression effects in other frameworks: scalar and fermion triplet leptogenesis, electroweak baryogenesis

In the above we have shown that a  $W_R$  discovered at current or future colliders would exclude any possibility to create a large enough baryon asymmetry from the decay of a  $N$ . However there exist other ways to induce successfully the baryon asymmetry through leptogenesis. In seesaw models this can be achieved from the decay of a scalar triplet to 2 leptons or from the decay of a fermion triplet to a lepton and Higgses, through diagrams involving another heavy state [13–15]. In these models there are washout effects from  $SU(2)_L$  interactions. These effects have been calculated in refs. [14, 15] and show that they are not large enough to rule-out leptogenesis even for masses as low as few TeV. For such low masses leptogenesis appears to be possible though only for asymmetries of order unity (i.e. assuming almost perfect resonance which requires e.g. large fine-tuning).

Since a  $W_R$  (or more generally any right-handed gauge boson) does not couple to left-handed triplets, its discovery at low scale would have no direct consequences for the triplet number density Boltzmann equation.

The discovery of a  $W_R$  at low scale would nevertheless provide a strong hint for the existence of  $N$ 's at low energy, see section 7. This would lead to 2 additional washout effects on the asymmetry produced by the triplet decays. First,  $\Delta L \neq 0$  scatterings involving both the  $W_R$  and  $N_R$ , figure 1, will be important (in the flavour channels coupling to the  $N$ 's) if both these particles have masses smaller or of order the triplet mass. Second, these  $N$ , through their Yukawa interactions, and together with sphalerons, could easily wash-out any previously produced lepton and baryon asymmetry, unless some of their Yukawa couplings are so suppressed that they preserve to a very good approximation at least one flavor number combination (which has not to be preserved in the triplet decay).

Putting all these effects together it can be checked that, the discovery of a  $W_R$  and a  $N$  would rule out the possibility to have any successful thermal leptogenesis from triplet decays at any scale as well, except for such kind of extreme flavour pattern.

Note that in the case of very low triplet mass a direct discovery of the triplets is possible through Drell-Yan pair production [31, 33].

Finally leptogenesis is also possible in more exotic models from the decay of  $SU(2)_{L,R}$  singlets, in case all the gauge interaction induced suppression effects considered in the above would be irrelevant for the decaying particle Boltzmann equation but still would be relevant for the  $Y_{\mathcal{L}}$  one. Similarly, electroweak baryogenesis with first order phase transition

from the presence of particles beyond the standard model around the electroweak scale, can be affected by the  $L$  violating interactions driven by a light  $W_R$  and/or light  $N$ , but could survive because these cannot erase the  $B$  asymmetry produced in this case. For electroweak baryogenesis at the right-handed scale [34] the effects could be large, and this would require a specific analysis.

## 7. $N$ and $W_R$ at colliders

We have shown this far to which (huge) extent the discovery of gauge interactions affecting the right-handed sector would cripple leptogenesis, offering - at least in the case of canonical neutrino decay leptogenesis - a rare opportunity of falsifying an otherwise particularly sturdy mechanism. This should provide additional motivation for this quest.

The discovery potential of LHC has been investigated for both massive right-handed neutrinos and gauge bosons associated to  $SU(2)_R$ ; in particular sensitivity plots corresponding to various stages of LHC operation can be found in [16, 17, 25], and scales of the order of 4 – 5 TeV in the best case are reached for  $W_R$ . Some attention should however be paid to the generality of the search. The "benchmark" just mentioned is reached under the assumption that at least one right-handed neutrino  $N$  is lighter than the  $W_R$ , and therefore that the process:  $p + p \rightarrow X + W_R \rightarrow X + N + l^-$  leads to an on-shell  $N$ , which can be reconstructed. Being a Majorana state, the  $N$  can decay indifferently into the channels  $l^- + u + \bar{d}$  or  $l^+ + \bar{u} + d$ , which, in connection with the production reaction leads to (non-resonant) dilepton signals of like or opposite charge in equal quantities. Same sign dilepton channels are particularly clean for background and its observation would establish the Majorana character of neutrino and  $N$  masses [35].

Given the importance for excluding leptogenesis, it may thus be worthwhile to go beyond this benchmark, and to examine the cases where either the  $W_R$ , the  $N$  or both are virtual.

The case of virtual  $N$  still gives a striking signature: namely, in equal amounts, 2 charged leptons of same or opposite sign + 2 jets, no missing energy, with the invariant mass resonating at  $m_{W_R}$ . The case of  $W_R$  heavier than the  $N$  is however of particular interest to us, even if the  $W_R$  only intervenes in a virtual way. In this case, the above process keeps the same overall signature, in particular equal amounts of like and opposite-sign dileptons, but resonance is only observed in the (lepton + 2 jets)- branch.

Only in the case where both  $N$  and  $W_R$  are both above threshold is the signature reduced to 2 jets + equal amounts of like or opposite charge dileptons.

It may also be worth pursuing other channels for detection of the  $W_R$ , in particular if the  $N$ 's are heavy. For this purpose, it is useful to note that, even if heavy  $N$ 's make the  $W_R$  leptonic decay impossible, it still couples to right-handed quarks whose mass is known. These quarks, being massive, also link to the left-handed sector. Hence the process  $p + p \rightarrow X + W_R^*$  followed by  $W_R^* \rightarrow t + \bar{b} \rightarrow \bar{b} + b + l^+ + \nu_L$ , the last decay occurring through an ordinary  $W_L$  ( $W_R^*$  stands here for either a real or a virtual  $W_R$ ) [36]. This possibility has been used at the Tevatron detectors [37] but not yet studied for LHC detectors. The interest in focusing on the top quark in the process is that it decays without having time

to hadronize, and therefore keeps the helicity correlations. In particular, the final lepton energy distribution is markedly softer [36] than in the similar process where both production and decay occur via  $W_L$ . A discovery through the top channel would not prove nevertheless that the  $W_R$  actually couples to the  $N$  but would be a strong hint for it.<sup>12</sup> We should finally mention the case where the right-handed neutrinos are (nearly) massless, in which case they cannot induce leptogenesis, but also cannot interfere with baryogenesis from another source. This case is difficult to characterize, as the right-handed closely resembles a heavier left-handed in most processes. Here again, the above-mentioned top quark intermediary channel, with its polarization effects would come to help.

## 8. Conclusion

We have shown that the discovery at LHC or future accelerators, of a  $W_R$  coupling to a right-handed neutrino and a right-handed charged lepton, would rule out the possibility to create any relevant lepton asymmetry from the decay of right-handed neutrinos, see figure 2. A  $W_R$  induces extra  $N$  decay channels inducing large dilution and washout effects, as well as very fast gauge scatterings (whose decoupling doesn't occur through Boltzmann suppression). We determined bounds on  $m_{W_R}$  and  $m_N$  for successful leptogenesis, given in figure 7 and eqs. (3.1) and (3.3). Similarly we discussed how the discovery of other particles generally expected in presence of right-handed gauge interactions, or of a  $Z'$ , could also affect leptogenesis, ruling it out too in some cases. Leptogenesis from the decay of scalar or fermion triplet would be also basically ruled out in presence of a  $N$  or both a  $N$  and a  $W_R$  around the TeV scale, unless there is a flavour symmetry to protect one flavour combination from the washout due to these states.

## Acknowledgments

The authors received partial support from the Belgian Science Policy (IAP VI-11), IISN, as well as from the NSF/PHY05-51164 grant. T.H. thanks the FNRS-FRS for support.

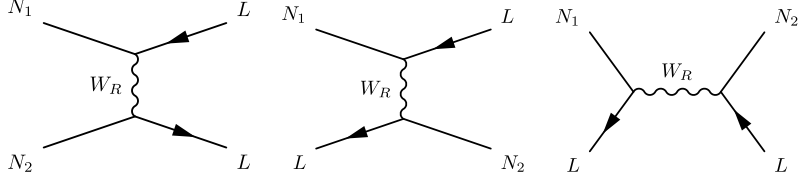
## A. Several right-handed neutrino case

With 2 right-handed neutrinos, and at the same level of approximation as for eqs. (2.10), (2.11),<sup>13</sup> we get the following Boltzmann equations:

$$\begin{aligned}
 zH(z)s Y'_{N_1} = & - \left( \frac{Y_{N_1}}{Y_{N_1}^{eq}} - 1 \right) \left( \gamma_{N_1}^{(l)} + \gamma_{N_1}^{(W_R)} + 2\gamma_{Hs}^{N_1} + 4\gamma_{Ht}^{N_1} + 2\gamma_{N_1u} + 2\gamma_{N_1d} + 2\gamma_{N_1e} \right) \\
 & - \left( \frac{Y_{N_1}^2}{Y_{N_1}^{eq2}} - 1 \right) \gamma_{N_1N_1}^{(W_R t)} - \left( \frac{Y_{N_1}Y_{N_2}}{Y_{N_1}^{eq}Y_{N_2}^{eq}} - 1 \right) \gamma_{N_1N_2}^{(W_R t)} \\
 & - \left( \frac{Y_{N_1}}{Y_{N_1}^{eq}} - \frac{Y_{N_2}}{Y_{N_2}^{eq}} \right) \left( \gamma_{N_1N_2}^{(W_R s)} + \gamma_{N_1N_2}^{(H,L)} \right)
 \end{aligned} \tag{A.1}$$

<sup>12</sup>Models where the  $W_R$  (or the  $Z'$ ) does not couple to the  $N$ , and therefore where it has little effect on leptogenesis, are with the  $SU(2)'_R (U(1)_N)$  subgroup of  $E_6$ , instead of the usual  $SU(2)_R$  [30].

<sup>13</sup>See footnote 4.



**Figure 9:** Scatterings involving 2  $N$ .

$$\begin{aligned}
zH(z)s Y'_{N_2} = & - \left( \frac{Y_{N_2}}{Y_{N_2}^{\text{eq}}} - 1 \right) \left( \gamma_{N_2}^{(l)} + \gamma_{N_2}^{(W_R)} + 2\gamma_{Hs}^{N_2} + 4\gamma_{Ht}^{N_2} + 2\gamma_{N_2u} + 2\gamma_{N_2d} + 2\gamma_{N_2e} \right) \\
& - \left( \frac{Y_{N_2}^2}{Y_{N_2}^{\text{eq}2}} - 1 \right) \gamma_{N_2N_2}^{(W_R t)} - \left( \frac{Y_{N_2}Y_{N_1}}{Y_{N_2}^{\text{eq}}Y_{N_1}^{\text{eq}}} - 1 \right) \gamma_{N_2N_1}^{(W_R t)} \\
& - \left( \frac{Y_{N_2}}{Y_{N_2}^{\text{eq}}} - \frac{Y_{N_1}}{Y_{N_1}^{\text{eq}}} \right) \left( \gamma_{N_2N_1}^{(W_R s)} + \gamma_{N_2N_1}^{(H,L)} \right) \tag{A.2}
\end{aligned}$$

$$\begin{aligned}
zH(z)s Y'_L = & \gamma_{N_1}^{(l)} \varepsilon_{N_1} \left( \frac{Y_{N_1}}{Y_{N_1}^{\text{eq}}} - 1 \right) + \gamma_{N_2}^{(l)} \varepsilon_{N_2} \left( \frac{Y_{N_2}}{Y_{N_2}^{\text{eq}}} - 1 \right) \\
& - \left( \gamma_{N_1}^{(l)} + \gamma_{N_1}^{(W_R)} + \gamma_{N_2}^{(l)} + \gamma_{N_2}^{(W_R)} \right) \frac{Y_L}{2Y_L^{\text{eq}}} \\
& - \frac{Y_L}{Y_L^{\text{eq}}} \left( 2\gamma_{N_s}^{\text{sub}} + 2\gamma_{Nt} + 2\gamma_{Ht}^{N_1} + 2\gamma_{Hs}^{N_1} \frac{Y_{N_1}}{Y_{N_1}^{\text{eq}}} + \gamma_{N_1u} + \gamma_{N_1d} + \gamma_{N_1e} \frac{Y_{N_1}}{Y_{N_1}^{\text{eq}}} \right. \\
& \left. + 2\gamma_{Ht}^{N_2} + 2\gamma_{Hs}^{N_2} \frac{Y_{N_2}}{Y_{N_2}^{\text{eq}}} + \gamma_{N_2u} + \gamma_{N_2d} + \gamma_{N_2e} \frac{Y_{N_2}}{Y_{N_2}^{\text{eq}}} \right) \tag{A.3}
\end{aligned}$$

$\gamma_{N_s}^{\text{sub}}$  and  $\gamma_{Nt}$  take into account the effects of the  $\Delta L = 2$  channels  $LH \leftrightarrow \bar{L}H$  and  $LL(\bar{L}\bar{L}) \leftrightarrow HH$  from both  $N_1$  and  $N_2$ .  $\gamma_{N_iN_j}^{(W_R t)}$  and  $\gamma_{N_iN_j}^{(W_R s)}$  parametrize the effects of the  $W_R$  mediated processes with 2 external  $N$ ,  $N_iN_j \leftrightarrow LL$  and  $N_iL \leftrightarrow N_jL$  respectively, as illustrated in figure 9. Similarly  $\gamma_{N_iN_j}^{(H,L)}$  parametrizes the effects of the Yukawa induced  $N_iL \leftrightarrow N_jL$  and  $N_iH \leftrightarrow N_jH$  scatterings mediated by a  $H$  and a  $L$  respectively. In these equations it is a very good approximation for the resonant case to take  $m_{N_1} = m_{N_2}$ ,  $\varepsilon_{N_1} = \varepsilon_{N_2}$ ,  $Y_{N_1}^{\text{eq}} = Y_{N_2}^{\text{eq}}$ ,  $\gamma_{N_1N_2}^{(H,L)} = \gamma_{N_2N_1}^{(H,L)}$ , as well as all gauge induced processes equal:  $\gamma_{N_1u,d,e} = \gamma_{N_2u,d,e}$ ,  $\gamma_{N_1N_2}^{(W_R t,s)} = \gamma_{N_2N_1}^{(W_R t,s)} = \gamma_{N_1N_1}^{(W_R t,s)} = \gamma_{N_2N_2}^{(W_R t,s)}$ .  $N_1$  and  $N_2$  can have significantly different effects only through their Yukawa coupling contributions.

To compare eqs. (2.10), (2.11) and eqs. (A.1), (A.2), (A.3) let us first note that the  $Y_{N_{1,2}}$  equations differ from the  $Y_N$  equation only through the  $\gamma_{N_iN_j}^{(W_R t,s)}$  and  $\gamma_{N_iN_j}^{(H,L)}$  terms. As in the one  $N$  case it can be checked that the  $\gamma_{N_iN_j}^{(W_R t)}$  terms have very little effects because their reaction rates are smaller than the  $\gamma_{Nu,d,e}$  ones (compare for example in figure 3a  $\gamma_{NN}$  with  $\gamma_{Ne} + \gamma_{Nu} + \gamma_{Nd}$ ). The  $\gamma_{N_iN_j}^{(W_R s)}$  terms on the other hand have a size similar to the one of  $\gamma_{Nu,d,s}$  but they are multiplied by  $Y_{N_2} - Y_{N_1}$ . This means that their effect is suppressed because those terms could be important only as long as the  $W_R$  effects ( $\gamma_{Nu,d,s}$  and  $\gamma_N^{(W_R)}$ ) dominate the thermalization of the  $N$ 's (with respect to the Yukawa induced processes), but these  $W_R$  effects equally affect  $Y_{N_1}$  and  $Y_{N_2}$ . Similarly it can be checked

that the  $\gamma_{N_i N_j}^{(HL\ s)}$  are of little importance. They are relevant only for very large values of both  $\tilde{m}_1$  and  $\tilde{m}_2$ , beyond the values of interest for our purpose. As a result all these terms can be neglected in eqs. (A.1), (A.2) and the evolution of  $Y_{N_1}$  and  $Y_{N_2}$  are essentially the same as the one of  $Y_N$  in eq. (2.10) replacing  $\tilde{m}$  by  $\tilde{m}_1$  and  $\tilde{m}_2$  respectively. There are no important differences at this level. Differences however can come from eq. (A.3) because this equation involves source and washout terms from both  $N_1$  and  $N_2$ . To discuss this equation it is useful to split it in two parts as follows

$$\begin{aligned}
 zH(z)s Y'_{\mathcal{L}a} &= \gamma_{N_1}^{(l)} \varepsilon_{N_1} \left( \frac{Y_{N_1}}{Y_{N_1}^{\text{eq}}} - 1 \right) - \left( \gamma_{N_1}^{(l)} + \gamma_{N_1}^{(WR)} + \gamma_{N_2}^{(l)} + \gamma_{N_2}^{(WR)} \right) \frac{Y_{\mathcal{L}a}}{2Y_L^{\text{eq}}} \\
 &\quad - \frac{Y_{\mathcal{L}a}}{Y_L^{\text{eq}}} \left( 2\gamma_{N_s}^{\text{sub}} + 2\gamma_{Nt} + 2\gamma_{Ht}^{N_1} + 2\gamma_{Hs}^{N_1} \frac{Y_{N_1}}{Y_{N_1}^{\text{eq}}} + \gamma_{N_1u} + \gamma_{N_1d} + \gamma_{N_1e} \frac{Y_{N_1}}{Y_{N_1}^{\text{eq}}} \right. \\
 &\quad \left. + 2\gamma_{Ht}^{N_2} + 2\gamma_{Hs}^{N_2} \frac{Y_{N_2}}{Y_{N_2}^{\text{eq}}} + \gamma_{N_2u} + \gamma_{N_2d} + \gamma_{N_2e} \frac{Y_{N_2}}{Y_{N_2}^{\text{eq}}} \right) \quad (\text{A.4})
 \end{aligned}$$

$$\begin{aligned}
 zH(z)s Y'_{\mathcal{L}b} &= \gamma_{N_2}^{(l)} \varepsilon_{N_2} \left( \frac{Y_{N_2}}{Y_{N_2}^{\text{eq}}} - 1 \right) - \left( \gamma_{N_1}^{(l)} + \gamma_{N_1}^{(WR)} + \gamma_{N_2}^{(l)} + \gamma_{N_2}^{(WR)} \right) \frac{Y_{\mathcal{L}b}}{2Y_L^{\text{eq}}} \\
 &\quad - \frac{Y_{\mathcal{L}b}}{Y_L^{\text{eq}}} \left( 2\gamma_{N_s}^{\text{sub}} + 2\gamma_{Nt} + 2\gamma_{Ht}^{N_1} + 2\gamma_{Hs}^{N_1} \frac{Y_{N_1}}{Y_{N_1}^{\text{eq}}} + \gamma_{N_1u} + \gamma_{N_1d} + \gamma_{N_1e} \frac{Y_{N_1}}{Y_{N_1}^{\text{eq}}} \right. \\
 &\quad \left. + 2\gamma_{Ht}^{N_2} + 2\gamma_{Hs}^{N_2} \frac{Y_{N_2}}{Y_{N_2}^{\text{eq}}} + \gamma_{N_2u} + \gamma_{N_2d} + \gamma_{N_2e} \frac{Y_{N_2}}{Y_{N_2}^{\text{eq}}} \right) \quad (\text{A.5})
 \end{aligned}$$

with  $Y_{\mathcal{L}} = Y_{\mathcal{L}a} + Y_{\mathcal{L}b}$ . Clearly comparing the  $Y_{\mathcal{L}a}$  ( $Y_{\mathcal{L}b}$ ) Boltzmann equations with the one  $N$  corresponding equation, eq. (2.11), one observes that these equations are the same except that eqs. (A.4), (A.5) involve additional washout terms from  $N_2$  ( $N_1$ ). Since these terms can only decrease<sup>14</sup> the absolute value of the lepton asymmetry obtained<sup>15</sup> one consequently gets

$$Y_{\mathcal{L}a}(m_N, \varepsilon_N, \tilde{m}_1, \tilde{m}_2) < Y_{\mathcal{L}}^{(1)}(m_N, \varepsilon_N, \tilde{m}_1) \quad (\text{A.6})$$

$$Y_{\mathcal{L}b}(m_N, \varepsilon_N, \tilde{m}_1, \tilde{m}_2) < Y_{\mathcal{L}}^{(1)}(m_N, \varepsilon_N, \tilde{m}_2) \quad (\text{A.7})$$

which gives

$$Y_{\mathcal{L}}(m_N, \varepsilon_N, \tilde{m}_1, \tilde{m}_2) < Y_{\mathcal{L}}^{(1)}(m_N, \varepsilon_N, \tilde{m}_1) + Y_{\mathcal{L}}^{(1)}(m_N, \varepsilon_N, \tilde{m}_2) \quad (\text{A.8})$$

with  $Y_{\mathcal{L}}^{(1)}$  which refers to the lepton number asymmetry obtained from eqs. (2.10), (2.11). This inequality has several consequences. (i) It means that if leptogenesis is ruled out in the one  $N$  case taking  $\varepsilon_N < 1$  (as above) it will be also ruled out in the 2  $N$  case if we take

<sup>14</sup>Except if in  $\gamma_{N_s}^{\text{sub}}$  and  $\gamma_{Nt}$  there is a destructive interference between the contribution of  $N_1$  and  $N_2$  but even so, from the effects of all other terms, the following inequalities hold (except for very large  $m_N$  close to  $10^{15}$  GeV which is not of interest for our purpose).

<sup>15</sup>Note that, due to the  $W_R$  effects it is a good approximation to start from thermal distributions of  $N_{1,2}$ , as explained above. Therefore there is no change of sign of  $Y_{\mathcal{L}}$  and the argument applies.

$\varepsilon_{N_{1,2}} < 1/2$  (which is the bound to be considered in this case, see ref. [15]). One just need to apply the results of figures 2 and 5 to both terms of eq. (A.8). (ii) As eq. (A.8) obviously also holds for the case where we neglect the  $W_R$  effects in the lepton number Boltzman equation, this conclusion remains true even if we play with flavour (applying to eq. (A.8) the results of figure 6). (iii) If, for a given value of  $m_N = m_{N_1} \simeq m_{N_2}$  and  $m_{W_R}$ , both  $\tilde{m}_1$  and  $\tilde{m}_2$  are outside the allowed range of  $\tilde{m}$  given in figure 7a, the lepton asymmetry produced will be too small. Numerically it can be checked also that this figure remains valid to a good approximation for the  $\tilde{m} = \tilde{m}_1 = \tilde{m}_2$  case. For  $m_{W_R}$  above  $\sim 50$  TeV the allowed region is shrunked by a hardly visible amount. As for the absolute lower bound on  $m_{W_R}$  it is larger in the 2  $N$  case than in the one  $N$  case (i.e. than the value 18 TeV above) but not by more than a few TeV. With more than 2 right-handed neutrinos these conclusions remain valid.

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