

## Affleck-Dine (pseudo)-Dirac neutrino genesis

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**ABSTRACT:** We consider the Affleck-Dine mechanism for leptogenesis in the minimal MSSM with Dirac or pseudo-Dirac neutrinos. The rolling of scalars along  $D$ -flat directions generates a left-right asymmetry in the sneutrino sector, only the left part of which is transferred to a baryon asymmetry via sphaleron transitions. In the pure Dirac case the baryon asymmetry of the Universe is thus mirrored by an equal and opposite asymmetry in the leptons. The mechanism is also found to work when the neutrinos are pseudo-Dirac. No additional field needs to be added to the MSSM other than the right-handed neutrino.

**KEYWORDS:** Cosmology of Theories beyond the SM, Baryogenesis.

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## Contents

<b>1. Introduction</b>	<b>1</b>
<b>2. The superpotential</b>	<b>2</b>
<b>3. The dynamics of the asymmetry</b>	<b>4</b>
<b>4. The baryon asymmetry</b>	<b>8</b>
<b>5. The pseudo-Dirac and mild see-saw cases</b>	<b>9</b>
<b>6. Conclusion</b>	<b>11</b>

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## 1. Introduction

It was noticed some time ago that leptogenesis, far from requiring lepton-number violating Majorana masses as in the original scenario [1], can in fact be implemented in the SM with purely Dirac neutrinos [2]. This mechanism, called *neutrinogenesis*, relies on the fact that  $(B + L)$ -violating transitions leave the right-handed sector unaffected [3]; as long as left-right equilibrating processes are small enough to be out of equilibrium, it is possible to 'hide' a right-handed lepton asymmetry from the sphaleron transitions. The idea of hiding lepton number in inert species has a long history [4] but works particularly effectively for the neutrinos; indeed it was shown in [2] that Dirac neutrinos easily satisfy this condition. The temporary left-handed lepton asymmetry can thus be processed before the electroweak phase transition into today's observed baryon asymmetry. It is only well after the phase transition that the neutrinos' Yukawa couplings come into equilibrium, by which time the sphalerons are quenched and the baryon asymmetry is locked in. The usefulness of this idea lies in the fact that Dirac neutrinos of the right size can arise in models where GUT scale degrees of freedom are integrated out because Yukawa couplings in such models can be naturally suppressed by factors of  $M_W/M_{GUT}$ . This possibility has received increased interest recently in the context of supergravity and effective models of string theories [5–10]. The fact that baryogenesis is also possible then leaves open the intriguing possibility that  $B - L$  is conserved in Nature or that neutrinos are in fact pseudo-Dirac rather than Majorana.

However the toy model used in [2] used an additional heavy Higgs-like doublet, because the scenario worked by 'drift and decay' as in original leptogenesis. In the present paper we point out that the Affleck-Dine (AD) mechanism [11] allows an extremely efficient implementation of neutrino genesis in just the MSSM with Dirac neutrinos. A  $(\tilde{\nu}_L - \tilde{\nu}_R)$  current can be produced through the rolling of scalars along their  $D$ -flat directions; although lepton-number is conserved, only left-handed lepton number can be converted to a baryon number through sphalerons, and the right-handed component is hidden by the smallness of the Yukawa coupling as before. The  $B + L$  number of the Universe is thus mirrored by an equal and opposite right-handed lepton number, until the right-handed (s)neutrino oscillations decay long after the electroweak phase transition. We should mention that AD neutrino genesis was proposed in ref. [12]. However in that work the AD field was considered to be an additional scalar field that was either Higgs-like, with  $SU(2)$  number, or a singlet appearing in higher order non-renormalizable interactions. The implementation here using only the  $D$ -flat directions of the MSSM itself can be thought of as the minimal realisation of AD neutrino genesis in the context of supersymmetry. Apart from testing whether AD can work in a  $B - L$  preserving MSSM this minimal scenario is naturally very predictive, because for example the CP violating potential arises from soft-supersymmetry breaking trilinear terms, the coupling of the AD field to the neutrinos is given by the neutrino mass, and so on.

We first discuss the evolution equations for the left-right ( $L - R$ ) asymmetry as a result of the Dirac mass term added to the MSSM superpotential. We will then evolve numerically the  $L - R$  asymmetry; using the equilibrium relations between  $B$  and  $L$ , this asymmetry will be converted to a baryon asymmetry. Finally we discuss the pseudo-Dirac and 'weak see-saw' cases. If neutrino genesis is to work a bound on the Yukawa couplings results from the requirement that the oscillations of the AD field decay after the electroweak phase transition. The nett result is a constraint on the size of the additional Majorana mass which turns out to be

$$M_R \lesssim 0.6 \left( \frac{0.05\text{eV}}{m_\nu} \right) \text{MeV}.$$

This bound is the strongest, being slightly more severe than the requirement that the baryon number generated be large enough. In short this means that the AD neutrino genesis is able to operate for all 'reasonable' Dirac and pseudo-Dirac scenarios in the neutrino mass sector. It even works for mild see-saw cases although the latter are probably excluded by nucleosynthesis.

## 2. The superpotential

Let us first introduce the right-handed neutrino superfield,  $\tilde{\mathbf{N}}$ , and add a Dirac mass term for the neutrinos in the superpotential:

$$\mathcal{V}_S \supset \lambda \mathbf{L}^i \epsilon_{ij} \mathbf{H}_u^j \tilde{\mathbf{N}}, \tag{2.1}$$

where  $\mathbf{L}$  is the left-handed lepton doublet and  $\mathbf{H}_u$  is the up-type Higgs. The gauge invariants  $\mathbf{LH}_u$  and  $\bar{\mathbf{N}}$  will be the important  $D$ -flat directions [13]; let us parameterise them as

$$\begin{aligned} L &= \frac{1}{\sqrt{2}} \begin{pmatrix} \phi \\ 0 \end{pmatrix}, \\ H_u &= \frac{1}{\sqrt{2}} \begin{pmatrix} 0 \\ \phi \end{pmatrix}, \\ \bar{N} &= \tilde{\nu} \end{aligned} \tag{2.2}$$

where non-bold letters stand for the scalar part of the superfields. The (SUSY-conserving) scalar potential arising from the added Dirac mass term is

$$\mathcal{V}_F = \left| \frac{\partial \mathcal{W}}{\partial L^a} \right|^2 + \left| \frac{\partial \mathcal{W}}{\partial H^b} \right|^2 + \left| \frac{\partial \mathcal{W}}{\partial \bar{\nu}} \right|^2 \tag{2.3}$$

$$= \frac{|\lambda|^2}{4} |\phi^2|^2 + |\lambda|^2 |\tilde{\nu}\phi|^2. \tag{2.4}$$

These tiny  $F$ -term contributions lift the  $\mathbf{LH}_u$  and  $\bar{\mathbf{N}}$  flat directions very slightly. In usual AD the directions considered would be both  $D$  and  $F$ -flat at the renormalizable level, and the flat direction would be lifted only by non-renormalizable (i.e. higher dimension) operators. Here the directions would be  $F$ -flat as well but for the neutrino mass. It is only because of the smallness of the latter that we can hope to send the field out to large enough VEVs to generate asymmetries. Note that conversely when we go on later to consider pseudo-Dirac neutrinos, then the scenario begins to run into difficulty if there is any significant see-saw effect at work in the neutrino mass matrices: a significant see-saw would imply larger Dirac Yukawa couplings and lift these directions more.

The AD mechanism of course requires additional CP violation, and here it comes from the soft-breaking sector;

$$\mathcal{V}_{SB} = m_\phi^2 |\phi|^2 + m_{\tilde{\nu}}^2 |\tilde{\nu}|^2 + (\lambda a \phi^2 \tilde{\nu} + h.c.). \tag{2.5}$$

As was discussed in [13, 14], soft-breaking terms also get a contribution from the non-zero Hubble constant in the early Universe and this is crucial as it drives the fields out to large values during inflation. We parameterise these as

$$\mathcal{V}_H = -c_\phi H^2 |\phi|^2 - c_\nu |\tilde{\nu}|^2 + (\lambda c_H H \phi^2 \tilde{\nu} + h.c.). \tag{2.6}$$

The overall potential for the scalar fields is thus

$$\begin{aligned} \mathcal{V} &= \mathcal{V}_F + \mathcal{V}_{SB} + \mathcal{V}_H \\ &= (m_\phi^2 - c_\phi H^2) |\phi|^2 + (m_{\tilde{\nu}}^2 - c_\nu H^2) |\tilde{\nu}|^2 + (\lambda(a + c_H H) \phi^2 \tilde{\nu} + h.c.) \\ &\quad + \frac{|\lambda|^2}{4} |\phi^2|^2 + |\lambda|^2 |\tilde{\nu}\phi|^2. \end{aligned} \tag{2.7}$$

For the flat directions to develop large expectation values during inflation, we need at least one of the fields to have a negative effective mass squared; here we will consider

$$(m_\phi^2 - c_\phi H^2) < 0. \tag{2.8}$$

Thus the Hubble induced terms in eq. (2.6) push the fields far from the origin. They are also important in introducing a time-dependence into the potential which guarantees that the AD mechanism will be in operation. To see this note that without these terms the potential always has a minimum where  $\text{sign}(\lambda a \phi^2 \bar{\nu}) = -1$  as this is the only trilinear term. Without an initial kick the fields will simply roll down this valley and no net lepton currents of any kind will be generated. The Hubble induced terms mean that the minimum is now at  $\text{sign}(\lambda(a + c_H H)\phi^2 \bar{\nu}) = -1$ . Thus even if  $c_H$  is real the phases of the fields have to become time dependent to track the instantaneous minimum: in effect the Hubble constant should kick the AD field for us wherever it starts out.

### 3. The dynamics of the asymmetry

To see these effects we proceed to examine how the asymmetry develops. First write the lepton number  $n_L$  as a sum of its right-handed and left-handed parts:

$$n_L = n_L^{(L)} + n_L^{(R)} \tag{3.1}$$

with  $n_L^{(L)}$  and  $n_L^{(R)}$  being in terms of our scalar fields

$$\begin{aligned} n_L^{(L)} &= \frac{i}{2} \left( \dot{\phi}^* \phi - \phi^* \dot{\phi} \right) \\ n_L^{(R)} &= -i \left( \dot{\bar{\nu}}^* \bar{\nu} - \bar{\nu}^* \dot{\bar{\nu}} \right) . \end{aligned} \tag{3.2}$$

The evolution equation for  $\phi$  is:

$$\ddot{\phi} + 3H\dot{\phi} + \frac{\partial \mathcal{V}}{\partial \phi^*} = 0 \tag{3.3}$$

and analogously for  $\bar{\nu}$ . Now using eq. (3.2) in eq. (3.3) and its conjugate, we find

$$\dot{n}_L^{(L)} + 3Hn_L^{(L)} = \text{Im} \left( \frac{\partial \mathcal{V}}{\partial \phi} \phi \right), \tag{3.4}$$

and again analogously for  $\bar{\nu}$ . From eq. (2.7) we see that the only imaginary terms are the  $a$ -terms and hence

$$\begin{aligned} \dot{n}_L^{(L)} + 3Hn_L^{(L)} &= 2\text{Im} (\lambda a \phi^2 \bar{\nu}) \\ \dot{n}_L^{(R)} + 3Hn_L^{(R)} &= -2\text{Im} (\lambda a \phi^2 \bar{\nu}) . \end{aligned} \tag{3.5}$$

We can see that lepton number eq. (3.1) is conserved,

$$\dot{n}_L + 3Hn_L = \frac{d}{dt} \left( n_L^{(L)} + n_L^{(R)} \right) + 3H \left( n_L^{(L)} + n_L^{(R)} \right) = 0, \tag{3.6}$$

but that the left-right asymmetry,  $n_L^{(L)} - n_L^{(R)} \equiv n_{LR}$ , has a non-trivial evolution:

$$\dot{n}_{LR} + 3Hn_{LR} = 4\text{Im} (\lambda a \phi^2 \bar{\nu}) . \tag{3.7}$$

The hope then is that this time dependence and CP asymmetry in the potential will generate a nett  $n_{LR}$ . If it does so this will feed through to the baryons via sphalerons. Before continuing we briefly establish the relation between  $n_L^{(R)}$  and the baryon number (before taking account of the effect of sphalerons,  $n_L^{(R)}$  is initially half the value of  $n_{LR}$  generated during the evolution of eq. (3.7)): the equilibrium ratio between lepton and baryon number under rapid sphaleron transitions was calculated in ref. [15, 16] for an SM like structure. In the present case we have an out-of-equilibrium right-handed Dirac neutrino in the analysis which simply holds a nett  $B - L$  and remains completely inert; therefore we can set  $(B - \hat{L}) = n_L^{(R)}$  where  $\hat{L}$  is the sum of all the leptons in equilibrium (i.e. excluding the right handed neutrino). In the MSSM we have an additional charged Higgs which changes the result from the SM; repeating the chemical potential analysis and assigning a chemical potential  $\mu_{B-\hat{L}}$  we find that above the electroweak phase transition for  $m$  Higgs fields

$$\begin{aligned} Y &= 16\mu_{B-\hat{L}} + (20 + 2m)\mu_Y \\ B - \hat{L} &= 26\mu_{B-\hat{L}} + 16\mu_Y \end{aligned} \tag{3.8}$$

while below it

$$\begin{aligned} Q &= 16\mu_{B-\hat{L}} + (44 + 4m)\mu_Y \\ B - \hat{L} &= 26\mu_{B-\hat{L}} + 16\mu_Q. \end{aligned} \tag{3.9}$$

Imposing  $(B - \hat{L}) = n_L^{(R)}$  this then translates into the following ratios:

$$\begin{aligned} B = L &= \frac{4(6 + m)}{66 + 13m} n_L^{(R)} = \frac{8}{23} n_L^{(R)} & T > T_{ew}, \\ B = L &= \frac{4(9 + m)}{111 + 13m} n_L^{(R)} = \frac{44}{137} n_L^{(R)} & T < T_{ew}, \end{aligned} \tag{3.10}$$

where  $L$  is the total lepton number including the right handed neutrinos and  $m = 2$  in the MSSM. (The only effect of the charged Higgs of the MSSM is to change the denominator of the last expression to  $111 + 13m$  rather than  $98 + 13m$ .) Note that in this pure Dirac case  $B - L$  is conserved and the final baryon number is approximately the LSP density, given by  $n_{DM} = \frac{23}{8} n_B$ , reminiscent of ideas pursued in a number of works [17–27]. This suggests the right handed sneutrino as the preferred LSP candidate since the LSP mass would then have to be of order 1GeV:

$$m_{DM} = \frac{8}{23} \frac{\Omega_{DM}}{\Omega_b} m_b. \tag{3.11}$$

Returning now to the dynamical evolution first note that the forcing term for  $n_{LR}$  evolution is time dependent and this is true generically when  $H \gg m_{3/2}$  even if the fields are sitting in the minimum because the minimum of the potential depends on  $H(t)$ . Indeed we have to minimise

$$\begin{aligned} \mathcal{V} &= (m_\phi^2 - c_\phi H^2) |\phi|^2 + (m_{\tilde{\nu}}^2 - c_\nu H^2) |\tilde{\nu}|^2 - 2|\lambda(a + c_H H)| |\phi|^2 |\tilde{\nu}| \\ &+ |\lambda|^2 |\phi^2|^2 + |\lambda|^2 |\tilde{\nu}\phi|^2. \end{aligned} \tag{3.12}$$

Taking the coefficient of  $|\bar{\nu}|$  positive, and for  $|c_\nu| H^2 \gg m_\nu^2$  and  $|c_\phi| H^2 \gg m_\phi^2$ , the minimum of the potential is given by

$$|\phi|_{\min}(t) \simeq \sqrt{\frac{c_\phi}{2}} \frac{H(t)}{\lambda} \tag{3.13}$$

$$|\bar{\nu}|_{\min}(t) \simeq \begin{cases} \frac{-c_\phi}{2c_\nu - c_\phi} \frac{|a|}{|\lambda|}, & c_H H \ll |a|, \\ \frac{-c_\phi}{2c_\nu - c_\phi} \frac{|H(t)|}{|\lambda|}, & c_H H \gg |a|. \end{cases}$$

One obvious difference between the present case and the more usual one in [13] is that here the initial field values have a lower bound of order  $\phi_{\min} \simeq a/\lambda \sim 10^{14}$  GeV even if  $H$  itself is much smaller than this value.

Let us now sketch the evolution from early after inflation to the electro-weak phase transition; we will confirm the picture with a numerical solution of the equations of motion as we go along. The numerically evolving  $n_{LR}$  is shown in fig.(1).

The mechanism requires that initially the fields are drawn far along the flat directions during inflation. This in turn requires the coefficient of  $|\phi|^2$  in eq. (2.7) to be negative as we have seen above. Inflation is then followed by an era of inflaton oscillation, during which the Universe is matter-dominated ( $H \sim 2/3t$ ). At this early stage the time dependence of the Hubble constant is still important in the evolution of the fields. Indeed both  $\phi$  and  $\bar{\nu}$  are following their evolution equations and, since  $H \gg m_\phi, m_{\bar{\nu}}$ , their motion is dominated by the falling value of the  $H^2$  Hubble induced mass-squared terms.

The behaviour of the fields in this phase is a major difference between the scenario we are considering here and the original AD scenario. The flat direction here is lifted by renormalizable terms; it is easy to see from the equations of motion that the distance of the fields from the instantaneous minimum drops as  $t^{-1}$ ; but eq. (3.13) tells us that the minimum itself drops as  $t^{-1}$  as well, so that during this matter dominated phase the fields are relatively undamped. We expect to see a long transient period and as our numerical analysis shows, this is indeed the case. To get a crude understanding of the behaviour in this phase, we can estimate the maximum amplitudes of the fields by assuming that the energy is constant in a co-moving volume:  $R^3 H^2 \phi_{\max}^2 = const$ . This gives  $\phi_{\max} = const$  which in turn gives  $n_{LR} = const$ . The detailed behaviour of  $n_{LR}$  is still rather complicated at this point; the fields are not yet executing regular cycles, and the current  $n_{LR}$  is rapidly flipping sign as the fields change their sense of rotation around the origin. This behaviour is in agreement with the arguments of refs.[13] where it was noted that only with nonrenormalizable terms of dimension 4 or higher do the AD fields follow the instantaneous minimum closely.

It is during matter-domination that the important  $H \sim m_{3/2}$  mark is reached. Below this point the Hubble induced terms in the effective potential become irrelevant to the evolution which is now dominated by the mass terms, and the behaviour changes markedly. The time dependence of the fields becomes more amenable to analytic approximation now since the equations of motion are nearly linear; indeed we may immediately use constancy of energy in a comoving volume argument to infer that  $R^3 m^2 \phi_{\max}^2 = const$  where  $m$  is the

mass of whichever field we are considering. If  $H = b/t$  this then suggests

$$\phi_{\max} \sim t^{-\frac{3b}{2}}, \quad (3.14)$$

which we indeed confirm numerically. In matter domination  $b = 2/3$  so that  $n_{LR}$  drops as  $t^{-1}$ . This can clearly be seen in fig.(1). Slightly later, but before reheating, the fields begin to exhibit “canonical” AD behaviour where they *do* execute regular cycles. Here we may approximate the real and imaginary components of fields as  $t^k \sin(mt)$ . Linearizing the equations of motion we find

$$\frac{k(k-1)\sin(m_\phi t)}{t^2} + \frac{2km_\phi \cos(m_\phi t)}{t} + \frac{3bm_\phi \cos(m_\phi t)}{t} = 0, \quad (3.15)$$

and then neglecting  $1/t^2$  terms we find  $k = -3b/2$  agreeing with the above.

The fields stay in this phase until reheating when the Universe becomes radiation dominated. We assume that reheating happens for  $T_R = 10^9$  GeV when the Hubble constant is  $H \sim T_R^2/M_{Pl} \sim 1$  GeV. We now have  $H = 1/2t$ , which gives  $k = -3/4$ , in accord with the late time behaviour of the Bessel function solutions of ref. [11]. The current then drops as  $n_{LR} \sim t^{-3/2}$  as regular matter until the electroweak phase transition. This behaviour is again confirmed by the numerical treatment: in particular fig.(1) shows the initial transients, the switch to AD behaviour, the  $t^{-1}$  and the  $t^{-3/2}$  behaviour.

It is in this final phase as the fields are rolling down to the minimum, that they capture a left-right asymmetry that is constant relative to the entropy. A positive  $n_{LR}$  means for instance that there is instantaneously more left-handed sneutrinos and right-handed anti-sneutrinos than left-handed anti-sneutrinos and right-handed sneutrinos, respectively: the left-handed sneutrinos are quickly turned into left-handed neutrinos through gaugino interactions. This can either go by decay with  $\Gamma \sim g_2^2 m_{\tilde{\nu}}$  or at high temperatures by scattering whose rate is

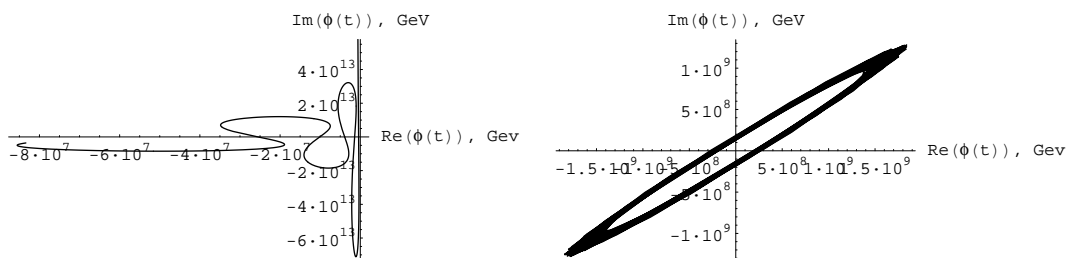
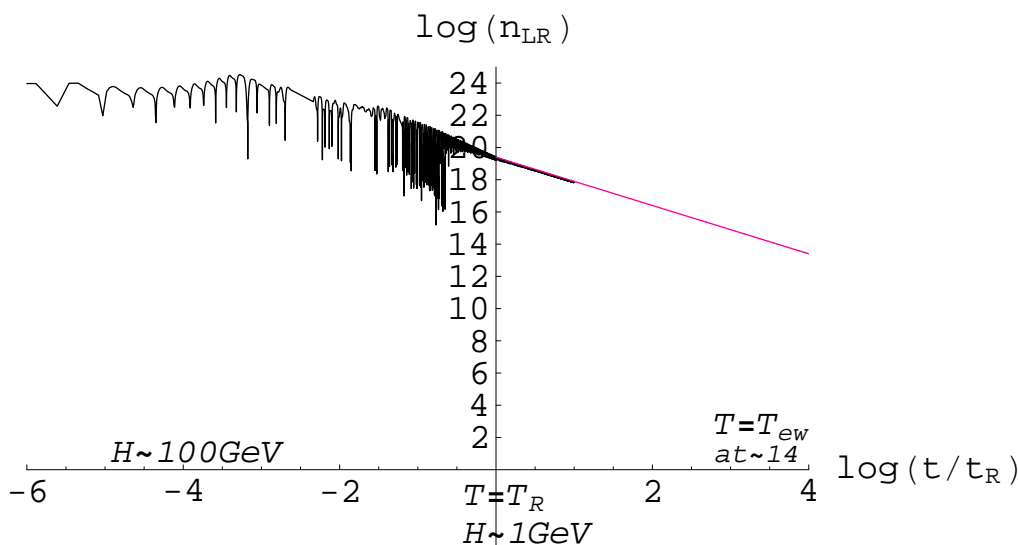
$$\Gamma \sim \frac{g_2^4}{m_{\tilde{W}, \tilde{B}}^4} T^5 \quad (3.16)$$

where the masses are understood to be thermal ones. All of the contributions are of the same order during the period we are considering when  $T \sim M_W$  and so sneutrino $\leftrightarrow$ neutrino conversion is in equilibrium. The sphaleron transitions transfer the left-handed neutrino asymmetry into a baryon asymmetry as described above. Above the electroweak phase transition this happens on a timescale of order  $\text{TeV}^{-1}$  which is essentially instantaneous; after the electroweak transition the sphalerons are switched off and the non-zero baryon number is frozen in [3, 28, 29]. Throughout, the right-handed (s)neutrinos remain inert until their oscillations decay through the coupling to the neutralino. We should therefore ensure that this happens at a temperature much lower than the electroweak transition temperature  $T_{ew}$ : the life-time of the sneutrino decay is

$$\tau_{\tilde{\nu}_R} \simeq \frac{4\pi}{\lambda^2} \frac{1}{m_{\tilde{\nu}}} B_{\text{Higgs}} \quad (3.17)$$

where  $B_{\text{Higgs}} < 1$  is the fraction of the neutralino that is made of the up Higgs. To be conservative we take  $B = 1$ , and find that  $\tau_{\tilde{\nu}_R} \gtrsim H^{-1}$  for all  $T \gtrsim \left(\frac{\lambda}{10^{-12}}\right) 100$  MeV. Note





**Figure 1:** Time evolution of the generated LR asymmetry. Parameters and initial conditions are as follows:  $m_\phi = 600$  GeV,  $m_{\tilde{\nu}} = 500$  GeV,  $a = e^{0.6i}100$  GeV,  $c_\phi = 1$ ,  $c_{\tilde{\nu}} = 0.8$ ,  $c_H = 0$ ,  $\lambda = 10^{-12}$ ,  $\phi(t_{in}) = \nu |\phi|_{\min}(t_{in})$ ,  $\tilde{\nu}(t_{in}) = |\tilde{\nu}|_{\min}(t_{in})$ ,  $\dot{\phi} = \dot{\tilde{\nu}} = 0$ , where the minima are given by the expressions in the text. The added line is matter evolution during radiation domination,  $t^{-3/2}$ . The behaviour of the  $\phi$  field is also shown for early (shortly before  $H \sim 100$ GeV) and late (post-reheating) times.

for later use that the mechanism stops working when

$$\lambda \gtrsim 10^{-9} \tag{3.18}$$

because the oscillations are damped before the electroweak phase transition takes place.

The discussion above of course assumes that the LSP is the usual neutralino. In the possibility we've mentioned above that the right-handed sneutrino itself be the LSP, the decay time for the sneutrino oscillations is even later.

#### 4. The baryon asymmetry

Having established the dynamical behaviour of the fields in some detail, let us now turn to the requirements to successfully generate a baryon number. First we have seen that

the oscillations in  $n_{LR}$  remain constant for an initial transient period when the inflaton oscillations are scaling like matter. They begin to scale like matter as well once the Hubble constant reaches  $\sim m_{3/2}$ . Therefore in order to evaluate the eventual  $n_{LR}$  asymmetry it is most useful to consider the relative densities when  $H \sim m_{3/2}$ ; the density in coherent inflaton oscillations is  $\rho_I \sim H^2 M_P^2 \sim m_{3/2}^2 M_P^2$ . At this stage the field VEVs are of order  $\phi, \bar{\nu} \sim |a/\lambda|$  as in eq. (3.13), so that the energy density in their oscillations is of order  $\rho_{\phi, \bar{\nu}} \sim m_{3/2}^2 |a/\lambda|^2$ . Since the latter also behaves like regular matter, we can use it to keep track of  $n_{LR}$  until the time of reheating:

$$\frac{\rho_{\phi, \bar{\nu}}}{\rho_I} \sim \frac{|a/\lambda|^2}{M_P^2}. \tag{4.1}$$

From reheating onwards it is the ratio with entropy that remains constant. Since  $\rho_{\phi, \bar{\nu}} = m_{\phi, \bar{\nu}} n_{\phi, \bar{\nu}}$  that is given by

$$\begin{aligned} \frac{n_{\phi, \bar{\nu}}}{s} &\approx \frac{|a/\lambda|^2 T_R}{M_P^2 m_\phi} \\ &= 10^{-9} \left| \frac{a}{100\text{GeV}} \right|^2 \left| \frac{10^{-12}}{\lambda} \right|^2 \left( \frac{T_R}{1\text{TeV}} \right) \left( \frac{100\text{GeV}}{m_\phi} \right). \end{aligned} \tag{4.2}$$

## 5. The pseudo-Dirac and mild see-saw cases

It is interesting to implement this scenario in the more general case where Majorana mass terms are included in the superpotential (2.1):

$$\mathcal{W} \supset \bar{\mathbf{N}} \lambda \mathbf{L}^a \epsilon_{ab} \mathbf{H}_u^b + M_R \bar{\mathbf{N}} \bar{\mathbf{N}} + \frac{M_L}{\langle h_u^0 \rangle^2} (\mathbf{L} \mathbf{H}_u)^2. \tag{5.1}$$

Such additional terms can arise in the same manner as the Dirac terms in the supergravity scenarios considered in ref. [5–10]; essentially the pure Dirac models require symmetries to prevent Majorana masses for the right-handed neutrinos that can be relaxed to allow non-renormalizable operators such as  $\mathbf{H}_u \mathbf{H}_d \bar{\mathbf{N}} \bar{\mathbf{N}} / M_{GUT}$ . For example such an operator could lead to left-handed Majorana masses  $M_L \sim 3 \times 10^{-5} - 7 \times 10^{-4}$  eV in the models considered in ref. [10]. In order to present as general a discussion as possible we will consider  $M_{L,R}$  to be arbitrary parameters and consider the question of *when AD neutrino genesis can work*. We will mainly focus on  $M_R$  since  $M_L \gtrsim \lambda v$  would give 6 active neutrinos which is certainly ruled out by nucleosynthesis; other than this we will consider  $M_{L,R}$  to be free parameters. Throughout the following discussion we shall assume that mass-squared differences given by measured neutrino oscillations are indicative of the actual masses.

The most immediate concern is how the new terms could affect the classical dynamics. Assuming for the moment that  $M_L = 0$ , the superpotential leads to the scalar potential (2.7):

$$\begin{aligned} \mathcal{V} &= (m_\phi^2 - c_\phi H^2) |\phi|^2 + (m_{\bar{\nu}}^2 - c_N H^2 + 4M_R^2) |\bar{\nu}|^2 \\ &\quad + (\lambda(a + c_H H) \phi^2 \bar{\nu} + \lambda M_R \phi^2 \bar{\nu}^* + \text{h.c.}) \\ &\quad + \frac{|\lambda|^2}{4} |\phi^2|^2 + |\lambda|^2 |\bar{\nu} \phi|^2. \end{aligned} \tag{5.2}$$

Clearly a new trilinear term has appeared which, following eq. (3.3), could affect the dynamics of the fields, and thus of the asymmetry if  $M_R \gtrsim a$ . However as we shall see such large values are not relevant for the AD scenario here.

Next, the lepton-number violating interactions introduced by the Majorana mass have to be constrained such that they do not erase the asymmetry. The rate of these interactions when  $M_R \neq 0$  is given by an exchange with  $\nu_R$  which for  $T \gg M_R$  is

$$\Gamma_{LV} \simeq \frac{\lambda^4 M_R^2}{T}, \quad (5.3)$$

Demanding that this rate be smaller than  $H$  for the duration of neutrinogenesis imposes

$$\lambda^4 M_R^2 \lesssim \frac{T_{ew}^3}{M_{Pl}} \quad (5.4)$$

which taking  $T_{ew} = 100\text{GeV}$  gives

$$\lambda^2 M_R \lesssim 3 \times 10^{-7} \text{GeV}. \quad (5.5)$$

As we shall see the resulting bounds are not very constraining. For  $M_L$  the lepton number violating exchanges are now suppressed only by the gauge couplings,  $g_2^4$ , rather than  $\lambda^4$ ; however this still gives an uninteresting bound  $M_L \lesssim 0.3\text{MeV}$ .

A different and more constraining bound comes from the Dirac Yukawas themselves; again considering  $M_L = 0$ , the light neutrino mass, given by

$$m_\nu = \frac{\sqrt{M_R^2 + (2\lambda v)^2} - M_R}{2}, \quad (5.6)$$

can require a  $\lambda$  substantially larger than  $10^{-12}$  due to see-saw effects once the Majorana contributions become dominant; any bound on  $\lambda$  coupled with the estimated neutrino mass puts an indirect bound on  $M_R$ . One such bound comes from the fact that if  $\lambda \gtrsim 10^{-8}$  then left- and right-handed neutrino can equilibrate above the electroweak phase transition, destroying any left-right asymmetry [2]. However as we saw a stricter bound,

$$\lambda \lesssim 10^{-9}, \quad (5.7)$$

comes from the requirement that the oscillations remain undamped until after the electroweak phase transition. In addition we of course require that the produced baryon number is sufficient; from eq. (4.2) we see that this is least constraining for maximum reheat temperature. Assuming that the gravitino reheat bound is satisfied  $T_R \lesssim 10^9\text{GeV}$  and that  $n_B/s \approx 10^{-10}$  we find the bound is (coincidentally) almost the same as in eq. (5.7),  $\lambda \lesssim 3 \times 10^{-9}$ . The added Majorana mass  $M_R$  must therefore respect both the constraints in (5.5) and (5.7). It is the latter which is most constraining; it gives

$$M_R \lesssim 0.6 \left( \frac{0.05\text{eV}}{m_\nu} \right) \text{MeV}. \quad (5.8)$$

This bound then prevents the see-saw mechanism from operating in all its glory, but a weak see-saw effect may still be present in the neutrinos while still allowing AD neutrinogenesis

to work with these flat directions. In addition note that the bound means that the evolution is virtually unaffected by the presence of the Majorana terms since  $M_R$  is indeed much smaller than both  $H$  and the trilinear term  $a$  driving the dynamics.

In this context, following eq. (4.2), there are two options for explaining why  $n_B/s \approx 10^{-10}$ . The first is the possibility that  $M_R \sim 1\text{MeV}$  and that there is a weak see-saw mechanism in operation. This seems rather unnatural since it introduces an additional mass-scale that itself requires explanation. Moreover, it has been argued that a 1MeV sterile neutrino with such a (relatively) large mixing angle ( $\sin^2 2\theta = 2 \times 10^{-7}$ ) would cause large amounts of sterile neutrino dark matter to be produced [30]. This scenario would then be forbidden by overclosure (and possibly other cosmological constraints — see [30]). In such a case, the remaining possibility is that the reheat temperature was of order 1TeV.<sup>1</sup>

## 6. Conclusion

In this paper we presented a minimal version of leptogenesis with Dirac sneutrinos in the MSSM, and showed that it can generate the observed baryon asymmetry of the Universe. No new fields need to be added to the MSSM apart from right-handed neutrinos. The mechanism works by first generating a  $\nu_L - \nu_R$  asymmetry using the AD mechanism, with  $D$ -flat directions involving sneutrinos and Higgses playing the role of the Affleck-Dine fields. The flat directions are appropriately lifted during inflation by the inclusion of finite-energy density SUSY-breaking terms which drives the VEVs to large values. As long as left-right equilibration is out of equilibrium before the electroweak phase transition (resulting in a bound on the Dirac neutrino Yukawas), the net left-handed lepton number can drive sphaleron transitions and ultimately create the observed baryon asymmetry. We also showed that the conditions on the smallness of the Yukawa couplings still allows the mechanism to be implemented for pseudo-Dirac neutrinos, and can in fact support a weak see-saw mechanism.

## Acknowledgments

The work of V.P. is funded by the Fonds Nature et Technologies du Québec. We thank Sacha Davidson for discussions at the beginning of this project. We are also grateful to K. Hamaguchi, R. Kitano, C. Lunardini and M. Ratz for comparison of results and for discussions, and to A.Mazumdar for useful comments.

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<sup>1</sup>It has been noted that such a low reheating temperature could arise in supersymmetry [31], if one of the  $D$  and  $F$ -flat directions acquires a VEV of order the Planck scale. (The flat directions required for the mechanism here do not get large enough VEVs to affect the reheat temperature by more than an order of magnitude).

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