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Dirac leptogenesis and anomalous $U(1)$

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Abstract. We consider Dirac leptogenesis in supersymmetric theories where the supersymmetry breaking is transmitted to the observable sector by an anomalous $U(1)$ symmetry. This kind of supersymmetry breaking is known to provide a solution to the μ problem and avoid large CP-violation effects. The asymmetries generated by the decays of heavy leptons do not suffer from wash-out due to the equilibration of left- and right-handed neutrinos thanks to the extreme smallness of the neutrino masses. The model ties up the smallness of the neutrino masses and the out-of-equilibrium nature of the heavy lepton decays with no tension with the overproduction of gravitinos.

The baryon asymmetry of the universe (BAU) is a quantity whose explanation lies on the interface between cosmology and particle physics. One of the attractive possibilities to explain this asymmetry is leptogenesis. Originally, the leptogenesis idea [1] relied upon the existence of heavy right-handed Majorana neutrinos whose decays into leptons generate baryon asymmetry of the universe via the fast sphaleron transitions. For this mechanism to work, the lepton number must be broken explicitly by the right-handed Majorana neutrinos[1]. Both baryon asymmetry of the universe and smallness of the neutrino masses are tied up to a common origin; namely right-handed heavy Majorana neutrinos.

Since the proposal of [1], an attractive alternative [3] view has arisen: Given that the smallness of the neutrino masses has been experimentally established, then there is no need to break the lepton number to have leptogenesis. Indeed, in a theory with exact lepton number conservation, a CP-violating decay of heavy lepton deposits equal asymmetries to left- and right-handed leptons, which are rapidly equilibrated via their Yukawa interactions. Neutrinos, however, constitute an exception to such fast processes. Indeed, due to the smallness of the neutrino masses, the equilibrium between the left- and right-handed neutrinos will not be attained until the temperature falls well below the weak scale. By this time, however, the asymmetry in the left-handed neutrinos will already be converted to BAU via the sphalerons. That this mechanism will produce enough baryon asymmetry can also be seen from a detailed analysis of the equilibration rate for the left- and the right-handed electron symmetries; these are much slower than those of all other species and much faster than those of neutrinos [4].

The main problem with this alternative view is that the mechanism which enables leptogenesis does not provide an explanation for the smallness of the neutrino masses. It is due to the empirical smallness of the neutrino masses; lepton numbers stored in left- and right-handed neutrinos will not equilibrate before the sphalerons go out of equilibrium [3].

In tackling the problem of generating small neutrino masses and enabling the leptogenesis, supersymmetric theories provide a viable arena. Indeed the very existence of the supersymmetry (SUSY) breaking scale can naturally generate small neutrino masses [5], not necessarily in the usual see-saw form.

In fact, recently it has been shown that the small neutrino masses and Dirac leptogenesis are compatible with each other [6], provided that the minimal model (MSSM) is extended by an extra $U(1)$ factor to forbid tree level Yukawa couplings for the neutrinos. The low energy particle spectrum consists of either an axion [7] or a Z' boson depending on whether this extra $U(1)$ invariance is global or local [8]. The reheat temperature after inflation must be small enough to avoid cosmological difficulties associated with gravitino overproduction, and, at the same time, must be high enough for the heavy mother leptons to decay out of equilibrium [9]. These cosmological constraints allow for small neutrino masses, if the neutrino Yukawa coupling is made hierarchically small.

As already pointed out in [6], for overcoming the difficulties associated with the gravitino problem, one possibility is to employ anomaly-mediated SUSY breaking in which the gravitino is sufficiently heavy [10]. On the other hand, the problem of tachyonic sleptons which is characteristic

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[1] The lepton number violation can arise also from the emission of right-handed neutrinos into bulk from the brane on which standard matter is localized [2].

to this scenario can be evaded, for instance, by incorporating a Fayet–Iliopoulos D-term associated with anomalous $U(1)$ [11]. Besides this, the anomalous $U(1)$ itself is known to mediate SUSY breaking either together with gravity [12,13], or together with gauge interactions [14]. Earlier works on SUSY breaking with anomalous $U(1)$ show that this scheme provides phenomenologically viable solutions to various hierarchy problems in the MSSM, namely the suppression of the supersymmetric CP-violation and the flavor-changing neutral current transitions, stabilization of the μ and B parameters of the MSSM Higgs potential to the weak scale, generation of appropriate Yukawa textures for explaining the hierarchy of the fermion masses, and finally, the small Dirac mass for neutrinos [12–19].

In general, there exist several $U(1)$ factors in effective theories arising from strings, and at least one factor is often anomalous (see e.g. [20]). The cancellation of the anomaly occurs via the Green–Schwarz mechanism [21], which requires both hidden and visible sector fields to be charged under the anomalous factor $U(1)_A$ thereby transmitting SUSY breaking from the former to the latter. As $U(1)_A$ is anomalous, $\mathrm{Tr}[Q] \neq 0$ ($\mathrm{Tr}[Q] = \sum_\alpha Q_\alpha$ where the Q_α are the $U(1)_A$ charges of the fields defined below) and a Fayet–Iliopoulos term $\xi \sim \mathcal{O}(M_{\mathrm{Pl}}^2)$ is generated [22]. It is this term which facilitates the SUSY breaking with soft terms $\mathcal{O}(\mathrm{TeV})$ [13].

In this work, we will discuss leptogenesis with Dirac neutrinos in theories where the SUSY breaking is transmitted to the observable sector by an anomalous gauge $U(1)$. Essentially what we are doing is to generalize the extra $U(1)$ symmetry of [6] to an anomalous one which facilitates the breaking of supersymmetry and generates the fermion Yukawa hierarchies via the $U(1)$ charges of the superfields. Once the predicted neutrino masses agree with the experiment with an appropriate assignment of the $U(1)_A$ charges, then decays of the heavy mother leptons are always out of equilibrium with no gravitino overproduction.

For definiteness we consider a pair of chiral superfields Φ_-, and Φ_+ which are neutral under the standard gauge group, but charged under $U(1)_A$ with respective charges -1 and $+1$. Then the $\mathcal{D}$-term contribution to the effective potential takes the form

$$\frac{g_A^2}{8}\,\mathcal{D}_A^2 = \frac{g_A^2}{8}\left(Q_\alpha\ \varphi_\alpha^\dagger\ \varphi_\alpha + |\Phi_+|^2 - |\Phi_-|^2 + \xi\right)^2 , \quad (1)$$

where g_A is the $U(1)_A$ gauge coupling, φ_α stands for the rest of the fields (the MSSM spectrum plus extra fields to be mentioned below), Q_α is the $U(1)_A$ charge of φ_α, and in string theories ξ is calculable:

$$\xi = \frac{g_A^2\ \mathrm{Tr}[Q]}{192\pi^2} M_{\mathrm{Pl}}^2 . \quad (2)$$

Note that ξ is completely determined for a given assignment of the $U(1)_A$ charges. The superpotential and the Kähler potential are invariant under the full gauge symmetry, and a possible candidate term for the superpotential is

$$W_0 = \mathcal{M}_\Phi\,\Phi_+\,\Phi_- , \quad (3)$$

where $\mathcal{M}_\Phi$ is calculable to be found to be $\mathcal{O}(\mathrm{TeV})$ if there exists a confining gauge theory at the intermediate scale [13, 23]. This mass term facilitates spontaneous breaking of SUSY together with $U(1)_A$. The minimization of the scalar potential yields $\langle\Phi_+\rangle = 0$, and $\langle\Phi_-\rangle = \sqrt{(\xi - 4\ \mathcal{M}_\Phi^2/g_A^2)}$, for the vacuum expectation values of the chiral superfields[2].

That $U(1)_A$ invariance is anomalous implies that each operator in the MSSM superpotential is dressed by appropriate powers of $\Phi_\pm/M_{\mathrm{Pl}}$ to achieve gauge invariance. Therefore, the superpotential of the model is essentially non-renormalizable and all interactions in the model, except for the gauge and soft-breaking couplings, are forbidden at the renormalizable level. This property generates the μ parameter and fermion mass textures and suppresses the electric dipole moments and flavor-changing neutral currents. Moreover, this model predicts naturally small Dirac masses [15, 18]. That a variety of hierarchy problems can be tackled succesfully stems from the $U(1)_A$ charges of the chiral fields which are predominantly positive [12] as also implied by the choice of positive ξ. This positiveness condition can be relaxed for several operators provided that $\mathrm{Tr}[Q] > 0$ persists. This is especially needed for the top quark [15], whose Yukawa interaction must appear at renormalizable level.

For realizing the Dirac leptogenesis, we extend the model above by two (heavy) lepton doublets $\mathcal{L}_-$, $\mathcal{L}_+$ and a right-handed neutrino N for each generation. These fields transform under $SU(2)_{\mathcal{L}} \times U(1)_Y \times U(1)_L \times U(1)_A$ as $N^i \sim (1,\ 0,\ -1,\ q_R^i)$, $\mathcal{L}_+^i \sim (2,\ -1/2,\ +1,\ q_+^i)$ and $\mathcal{L}_-^i \sim (2,\ 1/2,\ -1,\ q_-^i)$ where $U(1)_L$ stands for the lepton number. With this matter content, the following terms are added to the superpotential (3) of the model:

$$\begin{aligned} W_1 &= \lambda_+^{ij}\Phi_+\left(\frac{\Phi_-}{M_{\mathrm{Pl}}}\right)^{n_{\mathcal{L}_+}^{ij}} \mathcal{L}_+^i\cdot\mathcal{L}_-^j \\ &+ \lambda_-^{ij}\Phi_-\left(\frac{\Phi_-}{M_{\mathrm{Pl}}}\right)^{n_{\mathcal{L}}^{ij}} \mathcal{L}_+^i\cdot\mathcal{L}_-^j + h_+^{ij}\left(\frac{\Phi_-}{M_{\mathrm{Pl}}}\right)^{n_+^{ij}} N^i\mathcal{L}_+^j\cdot H_u \\ &+ h_-^{ij}\left(\frac{\Phi_-}{M_{\mathrm{Pl}}}\right)^{n_-^{ij}} L^i\cdot\mathcal{L}_-^j\Phi_+ + y^{ij}\left(\frac{\Phi_-}{M_{\mathrm{Pl}}}\right)^{n_L^{ij}} L^i\cdot H_u N^j\ , \end{aligned} \quad (4)$$

where i, j are the flavor indices, L^i is the lepton doublet of the ith generation, and H_u is the Higgs field. The matrices λ, h, y are in general complex and of order $\mathcal{O}(1)$ in size by naturalness. The first four operators in (4) define the masses and decay channels of heavy leptons after $U(1)_A$ breaking. One can check however that the contribution from the λ_+ term is negligible, as compared to the other

[2] Embedding of this model into supergravity does not spoil the breaking of supersymmetry but the ground state is modified by $\langle\Phi_+\rangle \to \sqrt{\xi/M_{\mathrm{Pl}}^2}\langle\Phi_-\rangle$. The light spectrum then contains the scalar Φ_+ with mass $M_+ \sim \mathcal{M}_\Phi$, and the gravitino with mass $m_{3/2} \sim \sqrt{\xi/M_{\mathrm{Pl}}^2}\ \mathcal{M}_\Phi$ [24].

three terms[3].

The exponents $n_{\mathcal{L}}, n_L, n_+$, and n_- are matrices in flavor space and they generate the textures. The last operator, for instance, generates the neutrino mass matrix [18]. The resulting neutrino masses will agree with experiment, if a typical diagonal entry of n_L is ~ 16 corresponding to $\xi \sim 10^{-2} M_{\rm Pl}^2$, which we will assume hereafter.

Gauge invariance imposes a total of 36 equations among the charges and matrices $n_{\mathcal{L}}$, n_L, n_+, and n_-. One immediate implication of gauge invariance is that

$$n_+^{\rm T} + n_- = n_{\mathcal{L}} + n_L + 2 \,, \qquad (5)$$

which forms a non-trivial constraint on the heavy lepton couplings to light fields N, L, H_u, and Φ_+. The mass matrix of the heavy leptons is given by

$$M_{\mathcal{L}} \sim \lambda_- \left(\frac{\xi}{M_{\rm Pl}^2}\right)^{(n_{\mathcal{L}}+1)/2} M_{\rm Pl} \,, \qquad (6)$$

where the flavor indices are suppressed. Similar to the kaon and Higgs systems with CP-violation, this mass matrix develops CP-violating entries via complex Yukawa matrices h_+, h_- through loop corrections. This indirect CP-violation effect originates from the mixing between the tree level decays $\mathcal{L}_+^i \to N^j H_u$, $\mathcal{L}_-^i \to L^j \Phi_+$, and the absorptive parts of the one loop wave function renormalization diagrams [25]. For definiteness we take the mass matrix of heavy leptons to be real and diagonal, $M_{\mathcal{L}} = {\rm diag}\,(M_1, M_2, M_3)$, with $M_1 \ll M_2 \ll M_3$. When M_1 is sufficiently small compared to $M_{2,3}$, the asymmetry is dominated by the decays of $\mathcal{L}_\pm^1$. For simplicity of the analysis, we consider only the lightest and next-to-lightest leptons neglecting the third (its presence provides more sources for generating asymmetry). For instance, the difference between N production and N^c production, summed over all final lepton generations, is given by [25]

$$\begin{aligned} \varepsilon_{N^c} &= \frac{\Gamma\left[\mathcal{L}_+^1 \to N^c H_u^c\right] - \Gamma\left[{\mathcal{L}_+^1}^c \to N H_u\right]}{\Gamma\left[\mathcal{L}_+^1\right]} \qquad (7) \\ &= \frac{1}{4\pi}\,\frac{M_2 M_1}{M_2^2 - M_1^2}\,\frac{{\rm Im}\left[\left(H_-^\dagger H_-\right)_{12}\left(H_+ H_+^\dagger\right)_{12}\right]}{\left(H_-^\dagger H_-\right)_{11} + \left(H_+^\dagger H_+\right)_{11}} \,, \end{aligned}$$

where we introduced the dressed Yukawa couplings

$$H_\pm = h_\pm \left(\frac{\xi}{M_{\rm Pl}^2}\right)^{n_\pm/2} \,, \qquad (8)$$

whose phases are generated by $h_\pm$ which are $\mathcal{O}(1)$ in size by naturalness. For the other asymmetries, a direct calculation gives $\varepsilon_L = -\varepsilon_{L^c} = \varepsilon_N = -\varepsilon_{N^c}$. These asymmetries enable the deposition of lepton number in left- and right-handed neutrinos and sneutrinos.

One of the main difficulties in a realistic leptogenesis scenario is to prevent the asymmetries generated above from wash-out. To evade this difficulty, the decay rate $\Gamma\left[\mathcal{L}^1\right]$ must be sufficiently small compared to the expansion rate of the universe $H(M_1)$. Indeed, the out-of-equilibrium conditions occur when the temperature falls below M_1 where the inverse processes are effectively blocked. This requires $\Gamma\left[\mathcal{L}^1\right] \lesssim 2H(M_1)$, in which case the abundance of heavy leptons increases, and this generates the requisite departure from the thermal equilibrium. In this limit, the heavy leptons drift and decay, and the asymmetry generated is converted into BAU with no suppression from the inverse decay processes.

Imposition of the out-of-equilibrium condition gives $n_{\rm o} \gtrsim (n_{\mathcal{L}}^{11} - 2)/2$ where $n_{\rm o} = {\rm Min}\,\{n_+^{11}, n_+^{12}, n_-^{11}, n_-^{21}\}$. If this condition is marginally satisfied, then one must explicitly solve the Boltzmann equations to see to what extent the decays are out of equilibrium. However, a quick glance at the condition (5) shows that this condition is well satisfied and the heavy lepton decays are far out of equilibrium provided that the neutrino masses are in the ballpark of the experimental results. Indeed, for $M_1 \sim 10^8\,$GeV, for instance, one has $n_0 \gtrsim 4$, whose actual value is at least three times larger than this bound once the condition (5) is taken into account. Therefore, the extreme lightness of the neutrinos guarantees that the heavy lepton doublets will decay well out of equilibrium.

Since the drift-and-decay limit well applies for heavy lepton decays, the stored lepton asymmetries are given by

$$\begin{aligned} L_R &\sim \frac{\varepsilon_N - \varepsilon_{N^c}}{g_*} \sim -2\frac{\varepsilon_{N^c}}{g_*}, \\ L_L &\sim \frac{\varepsilon_L - \varepsilon_{L^c}}{g_*} \sim -2\frac{\varepsilon_{N^c}}{g_*} \,, \end{aligned} \qquad (9)$$

where the sphalerons, which are in thermal equilibrium from $T \sim M_1$ down to the electroweak phase transition temperature, convert the asymmetry in left-handed neutrinos into the BAU [1]:

$$B = -\frac{28}{51} L_L \,. \qquad (10)$$

Due to the extreme smallness of the Higgs–ν_L–ν_R coupling compared to the gauge and Yukawa couplings, equilibration of L_R and L_L is not possible in time scales in which the sphaleron transitions are in equilibrium. Therefore, the asymmetry stored in the left-handed neutrinos will be converted to BAU [3, 4]. In addition to neutrinos, the asymmetries stored in sneutrinos can also contribute to BAU. However, in this case the equilibration processes are much faster depending on the size of the soft trilinear couplings. In case they do not equilibrate completely, they will contribute to BAU [26].

The produced baryon asymmetry (7) depends on the model parameters, in particular, on the Yukawa matrices $H_\pm$ and the masses of the heavy lepton doublets. For a

[3] Similar to neutrinos, all quark and charged lepton masses follow from non-renormalizable operators, and, therefore, their textures are generated via their $U(1)_A$ charges. The top quark being an exception, in general, all operators will have large charges under $U(1)_A$ to generate the small entries of the mass matrices [18]. Moreover, the suppression of the EDMs and FCNCs can both be accomplished via the charges of fields under anomalous $U(1)_A$ [15].

simple estimation, we assume that all entries of n_+ are similar in size to those of n_-; thus we take $n_+ \sim n_-$ entry-by-entry. Then the baryon asymmetry induced by the lepton asymmetry (7) depends on $h_\pm$, and the masses of the heavy leptons. Using (6), we get

$$B \sim 10^{-3} \left(\frac{\xi}{M_{\rm Pl}^2} \right)^{(n_{\mathcal{L}}^{11} - n_{\mathcal{L}}^{22})/2} \sin\phi \,, \qquad (11)$$

where ϕ is the relative phase between h_- and h_+, each of which is taken to be of order $\mathcal{O}(1)$ in size. The above expression produces enough asymmetry, with $\sin\phi \sim 1$, if $(n_{\mathcal{L}}^{11} - n_{\mathcal{L}}^{22}) \approx 7$. For instance, one may take $M_1 \sim 10^8$ GeV and $M_2 \sim 10^{15}$ GeV to generate sufficient BAU.

In SUSY models one of the tightest constraints come from the overproduction of gravitinos. Their decay products can disassociate the light elements and spoil the predictions of big-bang nucleosynthesis (BBN). There is, thus, always a connection between the reheat temperature (that must be high enough to allow for out-of-equilibrium decays of $\mathcal{L}^1$) and the mass of the gravitino [9]. For instance, $T_{\rm rh} \lesssim 10^7$ GeV, for $m_{3/2} \sim 100$ GeV and $T_{\rm rh} \lesssim 10^9$ GeV, for $m_{3/2} \sim 1$ TeV. In the model under consideration, $m_{3/2} \sim \sqrt{\xi/M_{\rm Pl}^2}\, M_\phi$ and it can be as low as 100 GeV for $M_\phi \sim 1$ TeV. Then, to prevent the overproduction of gravitinos, one can choose for instance, $n_{\mathcal{L}}^{11} \approx 11$ for having $M_1 \sim 10^6$ GeV. For higher values of the SUSY breaking scale M_ϕ (say ~ 10 TeV), the gravitino problem can be evaded with larger M_1 values. One notices that for smaller values of M_1, $n_{\mathcal{L}}^{11}$ becomes larger, and the heavy lepton decays fall farther out of equilibrium.

In the light of the above analysis we conclude that Dirac leptogenesis can generate sufficient BAU within the framework of SUSY breaking mediated by anomalous $U(1)$. The model is capable of explaining the small neutrino masses and leptogenesis consistently with no overproduction of gravitinos. The asymmetries generated by the heavy lepton decays cannot be washed out or erased via the equilibration of left- and right-handed neutrinos due to the extreme smallness of the neutrino masses.

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