Supersymmetric leptogenesis and the gravitino bound

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ABSTRACT

Supersymmetric thermal leptogenesis with a hierarchical right-handed neutrino mass spectrum requires the mass of the lightest right-handed neutrino to be heavier than about 10^9 GeV. This is in conflict with the upper bound on the reheating temperature which is found by imposing that the gravitinos generated during the reheating stage after inflation do not jeopardize successful nucleosynthesis. In this Letter we show that a solution to this tension is actually already incorporated in the framework, because of the presence of flat directions in the supersymmetric scalar potential. Massive right-handed neutrinos are efficiently produced non-thermally and the observed baryon asymmetry can be explained even for a reheating temperature respecting the gravitino bound if two conditions are satisfied: the initial value of the flat direction must be close to Planckian values and the phase-dependent terms in the flat direction potential are either vanishing or sufficiently small.

The observed baryon number asymmetry (normalized with respect to the entropy density) of the Universe $Y_B = (0.87 \pm 0.03) \times 10^{-10}$ [1] can be explained by the mechanism of thermal leptogenesis [2,3], the simplest implementation of this mechanism being realised by adding to the Standard Model (SM) three heavy right-handed (RH) neutrinos. In thermal leptogenesis the heavy RH neutrinos are produced by thermal scatterings after inflation and subsequently decay out-of-equilibrium in a lepton number and CP-violating way. The dynamically generated lepton asymmetry is then converted into a baryon asymmetry due to $(B + L)$-violating sphaleron interactions [4].

If RH neutrinos are hierarchical in mass, successful leptogenesis requires that the mass M_1 of the lightest RH neutrino N_1 is larger than 2×10^9 GeV, for vanishing initial N_1 density [5]. This lower limit on M_1 is reduced to 5×10^8 GeV when N_1 is initially in thermal equilibrium and to 2×10^7 GeV when N_1 initially dominates the energy density of the Universe [6]. These results do not substantially change when flavour effects are accounted for [7]. Hence, in the standard framework of thermal leptogenesis, the required reheating temperature after inflation $T_{\rm RH}$ cannot be lower than about 2×10^9 GeV [6]. In supersymmetric scenarios this is in conflict with the upper bound on the reheating temperature necessary to avoid the overproduction of gravitinos during reheating [8].

Being only gravitationally coupled to SM particles (and their supersymmetric partners), gravitinos decay very late jeopardizing the successful predictions of Big Bang nucleosynthesis. This does not happen, however, if gravitinos are not efficiently produced during reheating, that is if the reheating temperature $T_{\rm RH}$ is small enough. For gravitino masses in the natural range from 100 GeV to 1 TeV, within the minimal supergravity framework, the reheating temperature should be smaller than about 10^5–10^7 GeV [8], depending on the chosen values of the supersymmetric parameters and of the primordial element abundances.

The severe bound on the reheating temperature makes the thermal generation of the RH neutrinos impossible, thus rendering the supersymmetric thermal leptogenesis scenario not viable if RH neutrinos are hierarchical. Of course, there are several ways out to this drawback. First of all, one can modify the usual assumptions on gravitinos. If the gravitino is stable, the nucleosynthesis limit depends on the nature of the next-to-lightest supersymmetric particle, but values of $T_{\rm RH}$ even larger than 10^9 GeV can be obtained [9]. Assuming the existence of small R-parity violation, the next-to-lightest supersymmetric particle can decay before the onset of supersymmetry, evading the bound on $T_{\rm RH}$ [10]. Also, gravitinos lighter than 1 KeV (as possible in gauge mediation) or heavier than about 50 TeV (as possible in anomaly mediation) avoid the stringent limits on $T_{\rm RH}$. Alternatively, one can modify the standard mechanism of leptogenesis, and rely on supersymmetric resonant leptogenesis [11] or soft leptogenesis [12]. Indeed, in resonant leptogenesis the RH neutrinos are nearly degenerate in mass and self-energy contributions to the CP asymmetries are

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enhanced, thus producing the correct baryon asymmetry even at temperatures as low as the TeV. Soft leptogenesis can be successful for values of the mass M_1 of the lightest RH neutrino as low as 10^6 GeV. Another interesting variation is the case in which the right-handed sneutrino develops a large amplitude, dominating the total energy density [13]. Then the sneutrino decay reheats the universe, producing a lepton asymmetry, where values of T_{RH} as low as 10^6 GeV do not cause a gravitino problem. Finally, one can modify the standard thermal production mechanism of N_1. The lightest RH neutrinos can be produced non-thermally either during the preheating stage [14], or from the inflaton decays [15] or from quantum fluctuations [16].

In this Letter, we would like to show that a solution to the tension between supersymmetric leptogenesis with hierarchical RH neutrinos and the gravitino bound is in fact already rooted in one of the basic properties of the supersymmetric theory, that is the presence of flat directions in the scalar potential [17]. No new ingredient has to be added to the theory. Let us briefly sketch how the solution works. The F- and D-term flat directions are lifted because of the presence of the soft supersymmetry breaking terms in our vacuum, of possible non-renormalizable terms in the superpotential and of finite energy density terms in the potential proportional to the Hubble rate H during inflation [18]. As a consequence, the field ϕ along the flat direction will acquire a large vacuum expectation value (VEV). When, after inflation, the Hubble rate becomes of the order of the supersymmetry breaking mass $\tilde{m}$, the condensate starts oscillating around the true minimum of the potential which resides at $\phi = 0$. If the condensate passes close enough to the origin, the particles coupled to the condensate are efficiently created at the first passage. The produced particles become massive once the condensate continues its oscillation leaving the origin and may efficiently decay into other massive states, in our case RH neutrinos. The latter will subsequently decay to generate the final baryon asymmetry. The process allowing the generation of very massive states is called instant preheating [19] and represents a very efficient way of producing heavy states. In this sense, the solution we are proposing may be considered as a non-thermal production of RH neutrinos, but we stress that it does not involve any extra assumption such as a large coupling between the RH neutrinos and the inflaton field.

The generic potential for a supersymmetric flat direction ϕ is given by [18]

$$V(\phi) = (\tilde{m}^2 - cH^2)|\phi|^2 + \left(\lambda\frac{A + aH}{nM^{n-3}}\phi^n + \text{h.c.}\right) + |\lambda|^2\frac{|\phi|^{2n-2}}{M^{2n-6}},\quad (1)$$

where c, a and λ are constants of $\mathcal{O}(1)$, $\tilde{m}$ and A are the soft breaking mass terms of order the TeV scale, H is the Hubble rate, M is some large mass scale which we assume to be equal to the reduced Planck scale ($M = M_p = 2.4 \times 10^{18}$ GeV) and n is an integer larger than three. For $c > 0$ and $H \gg \tilde{m}$, the flat direction condensate acquires a VEV given by

$$|\phi_0| = \left(\frac{\beta H M^{n-3}}{\lambda}\right)^{1/(n-2)},\quad (2)$$

where β is a numerical constant which depends on a, c, and n. At the end of inflation, the inflaton field starts oscillating around the bottom of its potential and the Hubble rate decreases. As soon as $H \sim \tilde{m}/3$, the condensate starts rolling down towards its minimum at $\phi = 0$.

Now, if in the potential in Eq. (1) both terms proportional to A and aH are present and their relative phase $\theta_a - \theta_A$ does not vanish, the condensate $|\phi|e^{i\theta}$ will spiral around the origin at $\phi = 0$ with a nonvanishing $\dot{\theta}$ (possibly leading to a large baryon asymmetry through the Affleck–Dine mechanism [18,20]). In this case instant preheating does not occur and no heavy states are pro-

duced [21], unless several flat directions are simultaneously excited [22]. We will focus on the opposite case, when the condensate passes through the origin (or sufficiently close to it). This is easy to achieve without any fine-tuning [18] as it is enough to consider a flat direction which is lifted by a non-renormalizable superpotential term which contains a single field not in the flat direction and some number of fields which make up the flat direction [18],

$$W = \frac{\lambda}{M^{n-3}}\psi\phi^{n-1}.\quad (3)$$

For terms of this form, F_ψ is non-zero along the flat direction, but $W = 0$ along it. Examples of this type are represented by the direction ue which is lifted by $W = (\lambda/M)uude$, since $F_d^* = (\lambda/M)uue$ is non-zero along the direction, and by the Que direction which is lifted by the $n = 9$ superpotential $W = (\lambda/M)QuQuQuH_Dee$ since $F_{H_D}^* = (\lambda/M)QuQuQuH_Dee$ does not vanish [23]. If $W = 0$ along the flat direction, no phase-dependent terms are induced. Alternatively, the superpotential may vanish along the flat direction because of a discrete R-symmetry. In such a case, when W exactly vanishes, the potential during inflation has the form [18]

$$V(\phi) = H^2 M_p^2 f\left(|\phi|^2/M_p^2\right) + H^2 M_p^2 g\left(\phi^n/M_p^n\right),\quad (4)$$

and the typical initial value ϕ_0 for the condensate is $\mathcal{O}(M_p)$, rather than Eq. (2). For this reason we will treat ϕ_0 essentially as a free parameter in our analysis and not fixed by the relation (2). Finally we remark that the coefficients A and a depend on the specific form of the Kähler potential couplings and there are cases in which they are suppressed by inverse powers of M_p. For instance, if the inflaton is a composite field, it will appear in the Kähler potential only through bilinear combinations and $a \sim H/M_p$. In the case of D-term inflation [24] a vanishes identically and no phase-dependent terms are generated if along the flat direction $W = 0$.

From now on, we will consider a flat direction along which the induced A terms are suppressed and therefore the corresponding condensate will oscillate passing very close to the origin. Furthermore, we will focus on the flat direction involving the third generation quark u_3. When the condensate passes through the origin, it can efficiently produce states which are coupled to it. Let us consider the scalar Higgs H_U which is relevant for leptogenesis although, of course, other states will be produced as well. If the third generation is involved in the flat direction, the up-Higgs is coupled to the condensate through the Lagrangian term $h_t^2|\phi|^2|H_U|^2$. Its effective mass is therefore given by $m_{H_U}^2 = \tilde{m}_{H_U}^2 + h_t^2|\phi|^2$, where $\tilde{m}_{H_U}^2$ is the corresponding soft-breaking mass parameter. At the first passage through the origin, particle production takes place when adiabaticity is violated [19], $\dot{m}_{H_U}/m_{H_U}^2 \gtrsim 1$. This requires

$$\frac{|\dot{\phi}|}{h_t|\phi|^2} \sim \frac{\tilde{m}|\phi_0|}{h_t|\phi|^2} \gtrsim 1.\quad (5)$$

Up-Higgses can therefore be efficiently produced if $|\phi| \lesssim (\tilde{m}|\phi_0|/h_t)^{1/2} \equiv |\phi_*|$. As a result, particle production occurs nearly instantaneously, within a time

$$\Delta t_* \sim \frac{|\phi_*|}{|\dot{\phi}|} \sim \left(h_t\tilde{m}|\phi_0|\right)^{-1/2}.\quad (6)$$

The uncertainty principle implies that the created up-Higgses are generated with typical momentum [19]

$$k_* \sim \left(h_t\tilde{m}|\phi_0|\right)^{1/2}\quad (7)$$

and with a number density

$$n_{H_U} \sim \frac{k_*^3}{8\pi^3} \sim \frac{(h_t\tilde{m}|\phi_0|)^{3/2}}{8\pi^3}.\quad (8)$$

After the condensate has passed through the origin continuing its motion, the up-Higgses become heavier and heavier, having an effective mass $\sim h_t |\phi|$. When this mass becomes larger than the lightest RH neutrino mass M_1, the up-Higgses will promptly decay into the RH neutrinos N_1 (we suppose that the other RH neutrinos are much heavier than M_1) through the superpotential coupling $h_{ij} N_i \ell_j H_U$, where ℓ_j stands for the lepton doublet of flavour j and $i, j = 1, 2, 3$. Indeed, the H_U decay is prompt because the decay rate $\Gamma_D \sim \sum_j |h_{1j}|^2 h_t \phi/(8\pi)$ is faster than the oscillation rate $\Gamma_{osc} \sim \dot\phi/\phi$ as long as $\phi^2 > 8\pi \tilde{m}\phi_0/(\sum_j |h_{1j}|^2 h_t)$, which is certainly satisfied during the first oscillation. Moreover, if one of the h_{1j} is not too small, and Q_3 is not involved in the flat direction,[1] H_U will dominantly decay into $N_1\ell$, since any decay process occurring through top-Yukawa or gauge interaction is kinematically forbidden (or strongly suppressed) at large ϕ.

To estimate the maximum value M_1^{max} that can be generated we have to compute the maximum value ϕ^{max} achieved by the condensate during its first oscillation, after passing through the origin. The equation of motion for ϕ is

$$\ddot\phi + \tilde{m}^2 \phi = -h_t \frac{|\phi|}{\phi} n_{H_U}. \tag{9}$$

The term on the right-hand side corresponds to the ϕ-dependent energy density $m_{H_U}(\phi) n_{H_U}$ generated by the H_U particles produced when ϕ crosses the origin. It acts as a friction term damping the ϕ oscillations. Solving Eq. (9), we obtain

$$M_1^{max} \simeq h_t \phi^{max} = \frac{4\pi^3 \tilde{m}^{1/2} \phi_0^{1/2}}{h_t^{3/2}}$$
$$= 4 \times 10^{12} \text{ GeV} \left(\frac{\phi_0}{M_p}\right)^{1/2} \left(\frac{\tilde{m}}{100 \text{ GeV}}\right)^{1/2}, \tag{10}$$

where we have taken the top-Yukawa coupling $h_t \simeq 0.6$ at high-energy scales. Thus, very heavy RH neutrinos can be produced through this mechanism.

In first approximation, we can assume that all H_U decay into N_1 and the number density of the RH neutrinos is given by $n_{N_1} \sim n_{H_U} \sim (h_t \tilde{m}|\phi_0|)^{3/2}/8\pi^3$. When the mass of the up-Higgses decreases because the condensate, after reaching its maximum value at the first oscillation, starts decreasing again, the RH neutrinos may efficiently decay into up-Higgses and leptons and produce a lepton asymmetry $n_L \sim \epsilon n_{N_1}$ where the usual CP asymmetry ϵ is generated by the complex phases in the Yukawa couplings h_{ij}.

During all these stages, the inflaton field continues to oscillate around the minimum of its potential and will eventually decay into SM degrees of freedom giving rise to the reheating stage. Before reheating, the universe is matter dominated because of the inflaton oscillations and the scale factor increases as $a \sim H^{-2/3}$. The lepton asymmetry $n_L \sim \epsilon n_{N_1}$ produced during the first oscillation at $H_{osc} \sim \tilde{m}/3$ is diluted at the time of reheating by the factor $a_{osc}^3/a_{RH}^3 = H_{RH}^2/H_{osc}^2$. Expressing n_{N_1} through Eq. (8), we find that the baryon asymmetry $Y_B = (8/23)(n_L/s)(H_{RH}^2/H_{osc}^2)$ becomes

$$Y_B \sim \frac{9\epsilon h_t^{3/2} T_{RH}|\phi_0|^{3/2}}{92\pi^3 \tilde{m}^{1/2} M_p^2}$$
$$= 10^{-6}\epsilon \left(\frac{T_{RH}}{10^7 \text{ GeV}}\right) \left(\frac{|\phi_0|}{M_p}\right)^{3/2} \left(\frac{100 \text{ GeV}}{\tilde{m}}\right)^{1/2}. \tag{11}$$

Notice that in our estimate we have not inserted any wash-out factor. Indeed, as soon as the RH neutrinos decay, their energy density $\rho_{N_1} = M_1 n_{N_1}$ gets promptly converted into a "thermal" bath

with an effective temperature $\tilde{T} \sim (30\rho_{N_1}/g_*\pi^2)^{1/4}$ where g_* is the corresponding number of relativistic degrees of freedom. We estimate that $\tilde{T}$ is smaller than M_1 when

$$M_1 > 10^9 \text{ GeV} \left(\frac{|\phi_0|}{M_p}\right)^{1/2} \left(\frac{\tilde{m}}{100 \text{ GeV}}\right)^{1/2}. \tag{12}$$

As much heavier RH neutrinos are generated through the preheating stage, we may safely conclude that $\Delta L = 1$ inverse decays are not taking place. Similarly, one can show that the $\Delta L = 2$ processes are out-of-equilibrium. Finally, flavour effects [7] play no role in determining the final baryon asymmetry as $\Delta L = 1$ inverse decays are out-of-equilibrium. The maximum CP asymmetry parameter for normal hierarchical light neutrinos, in the supersymmetric case, is given by $\epsilon = 3M_1 m_3/(4\pi \langle H_U \rangle^2)$, where $m_3 = (\Delta m_{atm}^2)^{1/2}$ is the largest light neutrino mass. From Eq. (11), we therefore estimate that enough baryon asymmetry is generated if

$$M_1 \gtrsim 2 \times 10^{11} \text{ GeV} \left(\frac{10^7 \text{ GeV}}{T_{RH}}\right) \left(\frac{M_p}{|\phi_0|}\right)^{3/2} \left(\frac{\tilde{m}}{100 \text{ GeV}}\right)^{1/2}. \tag{13}$$

This limit, together with the result in Eq. (10), implies that a successful baryogenesis can occur only if $\phi_0 \gtrsim 0.2M_p(10^7 \text{ GeV}/T_{RH})^{1/2}$. The condensate of the flat direction has to start its oscillation from field values close to the reduced Planck mass. Notice that this limit on ϕ_0 is independent of h_t. However, the presence of the top Yukawa coupling is necessary to guarantee that the flat direction decays abundantly into H_U.

We conclude with some remarks. First, gravitinos are produced also during the instant preheating phase by scatterings of the quanta generated at the first oscillation of the condensate. It is easy to estimate that their abundance is $n_{3/2}/s \simeq 10^{-4}(T_{RH}/M_p)(\phi_0/M_p)^3$ and therefore it is never larger than the gravitino abundance produced at reheating by thermal scatterings, given by $n_{3/2}/s \simeq 2 \times 10^{-12}(T_{RH}/10^{10} \text{ GeV})$. Secondly, from Eq. (13) we infer that large values of the lightest RH neutrino mass M_1 are needed for the generation of a sufficiently large baryon asymmetry. However, we would like to point out that our mechanism can work also in models with smaller values of M_1, since the baryon asymmetry could be generated by the decays of the heavier RH neutrinos. Indeed, the up-Higgs may decay into the RH neutrinos N_2 (or N_3) instead into the lightest RH neutrino N_1 if the condensate reaches the value $\phi = \phi_{N_2} \equiv M_2/h_t$ before the up-Higgs decays into N_1's plus leptons. The time needed for the condensate to reach the value M_2/h_t is $\Delta t_{N_2} \sim \phi_{N_2}/\dot\phi \sim (M_2/h_t \tilde{m}\phi_0)$ and is smaller than the decay time of the up-Higgs into N_1's plus leptons if $\phi \lesssim (8\pi \tilde{m}\phi_0/\sum_j |h_{1j}|^2 M_2)$. Imposing that this critical value is larger than M_2/h_t, we find that the up-Higgs will promptly decay into N_2's rather than N_1's if

$$M_2 \lesssim \left(\frac{8\pi h_t \tilde{m}\phi_0}{\sum_j |h_{1j}|^2}\right)^{1/2}. \tag{14}$$

This condition can be satisfied if the Yukawas h_{ij} are hierarchical and $|h_{1j}| \ll 1$. If this is the case, one should replace M_1 with M_2 (or M_3) in Eqs. (12) and (13).

In conclusion, the observed baryon asymmetry can be explained within the supersymmetric leptogenesis scenarios for low reheating temperatures and a RH hierarchical mass spectrum, thus avoiding the gravitino bound, if two conditions are met: the initial value of the flat direction is close to Planckian values, and the phase-dependent terms in the flat direction potential are either vanishing or sufficiently small for the particle production to happen efficiently.

Acknowledgements

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[1] For the Que flat direction the $n = 9$ lifting superpotential contains Q_3 only if all the $n = 4$ lifting superpotentials $QQQL$, $QuQd$, $QuLe$ and $uude$ are present in the supersymmetric Lagrangian.

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