# Leptogenesis and rescattering in supersymmetric inflationary models

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The observed baryon asymmetry of the Universe can be due to the $B-L$ violating decay of heavy right handed (s)neutrinos. The amount of the asymmetry depends crucially on their number density. If the (s)neutrinos are generated thermally, in supersymmetric models there is limited parameter space leading to enough baryons. For this reason, several alternative mechanisms have been proposed. We discuss the nonperturbative production of sneutrino quanta by a direct coupling to the inflaton. This production dominates over the corresponding creation of neutrinos, and it can easily (i.e. even for a rather small inflaton-sneutrino coupling) lead to a sufficient baryon asymmetry. Light MSSM degrees of freedom are also significantly amplified via their coupling to the sneutrinos, during the rescattering phase which follows the nonperturbative production. This process, which mainly influences the (MSSM) $D-$flat directions, is very efficient as long as the sneutrinos quanta are in the relativistic regime. The rapid amplification of the light degrees of freedom may potentially lead to a gravitino problem. We estimate the gravitino production by means of a perturbative calculation, discussing the regime in which we expect it to be reliable.

The generation of the Baryon Asymmetry of the Universe (BAU) represents one of the puzzles of Cosmology. Three ingredients are required to achieve this task: baryon number violation, $C$ and $CP$ violation, and departure from thermal equilibrium. The baryon number violation can be challenging to implement, because it must be consistent with the current lower bound on the proton lifetime, $\tau_p \gtrsim 10^{32}$ years. The Standard Model (SM) is a $C$ and $CP$ violating theory, and contains non-perturbative $B+L$ violating interactions (sphalerons) which are rapid in the early Universe but unable to mediate proton decay. However, it seems difficult to use this baryon number violation to create

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the asymmetry in the SM and its more popular extensions. An attractive alternative is to generate a lepton asymmetry in some $C$, $CP$ and lepton number ($L$)-violating out-of-equilibrium interaction, and then allow the sphalerons to reprocess part of it into a baryon asymmetry. An appealing feature of this scenario is that while neutrino masses are experimentally observed (and are $L$-violating, if they are Majorana), there is still no evidence for baryon number violation.

The above idea is naturally implemented in the context of the see-saw, which is a minimal mechanism for generating neutrino masses much smaller than the ones of the charged leptons. Three right-handed (r.h.) neutrinos $N_i$ are added to the SM particle content, given Yukawa interactions with the lepton and Higgs doublets, and large Majorana masses. This gives the three light neutrinos very small masses, due to their small mixing with the heavy r.h. neutrinos through the Dirac mass. Grand Unified Theories (GUT) and their supersymmetric versions, that constitute natural candidates for the Physics beyond the SM, often contain r.h. neutrinos in their particle content. Here we consider the supersymmetric version of the see-saw mechanism, which is theoretically attractive because it addresses the hierarchy between the Higgs and r.h. neutrino masses. The r.h. (s)neutrinos of the see-saw can generate the BAU via leptogenesis, in a three steps process. First, some ($CP$ symmetric) number density of (s)neutrinos is created in the early Universe. Then, a lepton asymmetry is generated in their $CP$ violating out-of-equilibrium decay. Finally, the lepton asymmetry is partially reprocessed into a baryon one by the $B+L$ violating interactions, provided it is not washed out by lepton number violating scatterings. The following considerations are focused on the first step.

The most straightforward and cosmological model-independent mechanism to generate r.h. (s)neutrinos is via scattering in the thermal bath. However, the parameter space available is restricted in supergravity-motivated models. Indeed, unless some enhancement of the $CP$ asymmetry characterizing the r.h. (s)neutrino decay is present (which occurs for example if they are nearly degenerate in mass) the generation of a sufficient lepton number poses a rather strong lower bound on their mass. The final baryon asymmetry is indeed found to be

$$Y_B \simeq 10^{-10} \left(\frac{M_1}{10^{10}\,\mathrm{GeV}}\right) \left(\frac{m_3}{0.05\,\mathrm{eV}}\right) \left(\frac{\kappa}{0.1}\right) \delta_{\mathrm{CP}} \;, \tag{1}$$

where $M_1$ is the mass of the lightest r.h. (s)neutrino, $m_3$ the mass of the heaviest l.h. neutrino, $\kappa \simeq 10^{-2}$–$10^{-1}$ comes from the wash out of the asymmetry due to lepton number violating scattering (typically still in equilibrium when the sneutrino decays), and finally $|\delta_{\mathrm{CP}}| \leq 1$. For thermal production of the r.h. (s)neutrinos, this expression translates into a lower bound on the reheating temperature of the thermal bath, $T_{RH} \gtrsim M_1$. On the other hand, if supergravity is assumed, $T_{RH}$ must be below $10^9 - 10^{10}$ GeV, to avoid overproduction of gravitinos. The two requirements are compatible only provided a nearly maximal $CP$ asymmetry is present in the r.h. (s)neutrino decay, $|\delta_{\mathrm{CP}}| \simeq 1$. If this is not the case, alternative mechanisms of production for the r.h. (s)neutrino have to be considered.

Leptogenesis can be achieved if at least one r.h. sneutrino has a smaller mass than the Hubble parameter, *i.e.* [3] $M_N < H$, during inflation. In this case, quantum fluctuations of

[3]We will use $N$ as a shorthand for the superfield, its scalar and fermionic component. We explicitly

this sneutrino component are produced during the inflationary expansion, and amplified to generate a classical condensate. The decay of the condensate eventually generates the required lepton asymmetry. The above requirement $M_N < H$ is not trivially satisfied in a supergravity context, since supergravity corrections typically provide a mass precisely of order $H$ to any scalar field of the model. In this case, however, a suitable choice of the Kähler potential can induce a negative mass term $m^2_{\rm ind} \simeq -H^2$, so that a large expectation value will be generated for the sneutrino component during inflation. This also leads to the formation of a condensate during inflation, and to successful leptogenesis.

Large variances can be produced during inflation if the sneutrinos are not too strongly coupled to the inflaton field $\phi$, since this would generate a high effective mass which could fix $\langle \tilde{N} \rangle = 0$ during inflation. However, if one of the r.h. (s)neutrinos is coupled to the inflaton, there is the obvious possibility that a sufficient amount of (s)neutrino quanta is generated when the inflaton decays. Quite remarkably, for a rather wide range of models this decay occurs in a nonperturbative way (this is known in the literature as *preheating*). In models of chaotic inflation, this is due to the coherent oscillations of the inflaton field, which can be responsible for a parametric amplification of the bosonic fields to which the inflaton is coupled. It is important to remark that this resonant amplification does not require very high couplings between the inflaton and the produced fields. For a coupling of the form $(g^2/2)\,\phi^2\,N^2$ in the scalar potential, resonant amplification of the field $N$ already occurs for $g^2 \gtrsim 10^{-8}$, if the mass of $N$ is negligible at the end of inflation, and if a massive inflaton is considered. For a massless inflaton ($V(\phi) = \lambda\,\phi^4/4$), an efficient resonance is present also for much smaller values of $g$, since in this case the resonance is not halted by the expansion of the Universe.

If the produced particle is very massive ($M_N \gtrsim m_\phi$), the effectiveness of the resonance becomes a highly model dependent issue. Consider for example a model characterized by the potential $V(\phi) + (M_N + g\,\phi)^2\; N^2/2$. In this case, due to the high initial amplitude of the inflaton oscillations, the total mass of $N$ can vanish at some discrete points even for a coupling as small as $g^2 \sim 10^{-10}\;(M_N/m_\phi)^2$. Whenever $M_N + g\,\phi = 0$, parametric amplification of $N$ occurs. Although this choice of the potential may seem *ad hoc*, we note that it is the one which arises in supersymmetric models if both the r.h. sneutrino mass and interaction with the inflaton are encoded in the superpotential, $W(N) \supset M_N\,N^2 + g\,\phi\,N^2$. We regard this as a very natural possibility.

The nonperturbative production of r.h. neutrinos, with a mass term of the form $(M_N + g\,\phi)\,\bar{N}N$ has been recently discussed. It has been shown that a sufficient lepton asymmetry is generated if the mass of the r.h. neutrinos is higher than about $10^{14}$ GeV, and if their coupling to the inflaton satisfies $g \gtrsim 0.03$. Here we note that this high coupling can in principle destabilize through quantum effects the required flatness of the inflaton potential. This, in addition to the strong hierarchy between the r.h. neutrino mass and the electroweak scale, motivates the study of the supersymmetrized version of this mechanism.

In the supersymmetric version, the nonperturbative production of the r.h. sneutrinos is much more efficient than the one of the neutrinos. Due to supersymmetry, the inflaton couples with the same strength both to the r.h. neutrinos and to the sneutrinos, so that if

refer to the particle type when this could cause a confusion.

the former are produced at preheating this will also occur for the latter. However, while production of fermions is limited by Pauli blocking, the production of scalar particles at preheating is characterized by very large occupation numbers. The production of scalar quanta is so efficient that it typically leads to a strong amplification of all the other fields to which they are coupled, irrespective of the fact that the latter are directly coupled or not to the inflaton. This is a very turbulent process, dominated by the nonlinear effects (generically denoted as *rescattering*) caused by the very high occupation numbers of the fields involved. As a result, all these mutually interacting fields are left with highly excited spectra far from thermal equilibrium.

Rescattering strongly affects some of the outcomes of the analytical studies of preheating of bosons, which do not take into account all these nonlinear effects. However, due to their high occupation numbers, the scalar distribution show a classical (deterministic) evolution, which can be studied by means of numerical simulations on the lattice. Full numerical calculations on the lattice are rather extensive. We have found that the necessary computing time is reduced in the conformal case, that is with the inflaton potential $V(\phi) = \lambda\, \phi^4/4$ , and with a r.h. sneutrino mass which is negligible during the early stages of preheating. For this reason, in the computations we discuss below we fix $M_N = 10^{11}$ GeV, which is smaller than the Hubble parameter during inflation, but still high enough to require a nonthermal production of the sneutrinos. The numerical results show a very efficient production of r.h. sneutrinos and inflaton quanta at preheating/rescattering. Even for a coupling inflaton-sneutrino as small as $g^2 \sim \text{few} \times 10^{-12}$ , the produced quanta come to dominate the energy density of the Universe already within about the first 5 e-folds after the end of inflation. In particular, the energy density stored in sneutrinos is typically found to be a fraction of order one of the total energy density, so that a sufficient leptogenesis is easily achieved at their decay.

R.h. sneutrinos are coupled to Higgs fields and left handed (l.h.) leptons through the superpotential term $h\, N\, H\, L \subset W$ (responsible for the Yukawa interaction which provides a Dirac mass to the neutrinos). Thus, one may expect that quanta of the latter fields are amplified by the rescattering of the r.h. sneutrinos produced at preheating. Numerical simulations indeed show that the amplification occurs for a wide range of values of the coupling $h$ . An important role in determining how significant this effect can be is played by the mass of the r.h. sneutrinos. More precisely, when the sneutrino quanta become non relativistic (let us denote by $\hat{\eta}$ the time at which this happens) their rescattering effects become much less efficient. Thus, a strong amplification of the MSSM fields at rescattering can take place only if the coupling $h$ is sufficiently large so that the amplification occurs before $\hat{\eta}$ . As a consequence, for massive sneutrinos and for small values of $h$ , the number of MSSM quanta produced at rescattering is an increasing function of $h$ . However, the production is actually *disfavored* when the coupling $h$ becomes too high. This is simply due to energy conservation, since the energy associated to the interaction term between the sneutrino and the MSSM fields cannot be higher than the energy initially present in the sneutrino distribution (equivalently, one can say that, for a too high coupling $h$ , the non vanishing value of the sneutrinos gives a too high effective mass to the MSSM fields, which prevents them from being too strongly amplified).

An important remark is in order. When we speak about the amplification of MSSM fields coupled to the r.h. sneutrinos we have actually in mind amplification of $D-$flat directions (let us generally denote them by $X$). Indeed, $D-$terms provide a potential term of the form $\Delta V \sim g_G^2 |Y|^4$ for any scalar non flat direction $Y$ . Since $g_G$ is a gauge coupling ($g_G = O\left(10^{-1}\right)$), we expect such terms to prevent a strong amplification of $Y$ , again from energy conservation arguments. Another important issue which emerges when gauge interactions are considered is whether gauge fields themselves are amplified at rescattering. We believe that, at least in the model we are considering, also the amplification of gauge fields will be rather suppressed. The scalar distributions amplified at rescattering break much of the gauge symmetry of the model. This gives the corresponding gauge fields an effective mass in their dispersion relation (analogous to the thermal mass acquired by fields in a thermal bath) of the order $m^2 \sim g_G^2 \langle X^2 \rangle$ . As we extensively discuss in the paper, in the class of models we are considering the nonthermal distributions formed at rescattering are characterized by a typical momentum several orders of magnitude smaller than this mass scale. For this reason, one can expect that such heavy gauge fields cannot be strongly amplified. [4] In our opinion, an explicit check of these conjectures by means of numerical simulations could be of great interest, especially considering the great importance that gauge fields could have for the thermalization of the scalar distributions.

To conclude, we discuss the production of gravitino quanta from the scalar distributions generated at rescattering. We already mentioned that in order to avoid a thermal overproduction of gravitinos a lower bound has to be set on the reheating temperature $T_{RH}$ of the thermal bath. The requirement of a low reheating temperature can be seen as the demand that the inflaton decays sufficiently late, so that particles in the thermal bath have sufficiently low number densities and energies when they form. If $H \simeq 10^{12}$ GeV at the end of inflation, and if the scale factor $a$ is normalized to one at this time, the generation of the thermal bath cannot occur before $a \simeq 10^7$ . Gravitino overproduction is avoided by the fact that in the earlier times most of the energy density of the Universe is still stored in the coherent inflaton oscillations. On the contrary, we have already remarked that preheating/rescattering lead to a quick depletion of the zero mode in the first few e-folds after the end of inflation.

The question whether also the distributions formed at rescattering may lead to a gravitino problem is thus a very natural one. [5] To provide at least a partial answer to this question, we distinguish the period during which rescattering is actually effective from the successive longer thermalization era. The computation of the amount of gravitinos produced during the earlier stages of rescattering appears as a very difficult task. The numerical simulations valid in the case of bosonic fields indicate that a perturbative computation (with dominant $2 \to 2$ scatterings taken into account) can hardly reproduce the numerical results, and that probably $N \to 2$ processes ($N > 2$) have also to be taken into account. It is expected that the same problem will arise also for the computations of the quanta of gravitinos produced by the scalar distributions which are being forming at this stage. The end of rescattering/beginning of the thermalization period is instead characterized by a much slower evolution of the scalar distributions. In particular, the total occupation number of all the scalar fields is (approximately) conserved, which is interpreted by the fact that $2 \to 2$ processes are now determining the evolution of their

[4] Gauge fields which are not coupled to the fields generated at rescattering will not acquire this high effective mass. However, being uncoupled, they will not be amplified either.

[5] We acknowledge very useful discussions with Patrick B. Greene and Lev Kofman on this issue.

distributions. Motivated by this observation, we assume that $2 \rightarrow 2$ interactions are also the main source of production for gravitinos from this stage on.

In the thermal case, the gravitino production is dominated by processes having a gravitationally suppressed vertex (from which the gravitino is emitted) and a second vertex characterized by a gauge interaction with one outgoing gaugino. However, we believe that in the present context these interactions will be kinematically forbidden, due to the high effective mass-squared that gauginos acquire from their interaction with the scalar distributions (the argument follows the one already given for gauge fields). Once again we notice that the system is still effectively behaving as a condensate: the number densities of the scalar distributions are set by the quantity $\sqrt{\langle X^2 \rangle}$, which is much higher than the typical momenta of the distributions themselves. This generates a high effective mass for all the particles "strongly" coupled to these scalar fields. A further comparison with the case of a thermal distribution may be useful: in the latter case both the typical momenta and the effective masses are set by the only energy scale present, namely the temperature of the system. As should be clear from the above discussion, the thermalization of the distributions produced at rescattering necessarily proceeds through particle fusion. Only after a sufficiently prolonged stage of thermalization, will the system be sufficiently close to thermodynamical equilibrium so to render processes as the one discussed above kinematically allowed.

Numerical simulations show that if this class of processes is indeed kinematically suppressed, the production of gravitinos from the distributions formed at rescattering is sufficiently small. However, we remark that this analysis still leaves out the gravitino production which may have occurred at the earlier stages of the rescattering period. Whether this production may be sufficiently strong to overcome the limits from nucleosynthesis remains an open problem.

The results summarized in this talk, are discussed in details in ref. [1]. An extensive list of relevant bibliography can be found there.

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## References

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