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Baryogenesis from 'electrogenesis' in a scalar field dominated epoch

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Abstract: Scalar fields can play a dominant role in the dynamics of the Universe until shortly before nucleosynthesis. Examples are provided by domination by a kinetic mode of a scalar field, which may be both the inflaton and the late time 'quintessence', and also by more conventional models of reheating. The resultant modification to the pre-nucleosynthesis expansion rate can allow solely an asymmetry in right handed electrons to produce a net baryon asymmetry when reprocessed by the anomalous $B + L$ violating processes of the standard model. The production of such a source asymmetry — what we term 'electrogenesis' — requires no additional total B or L violation beyond that in the standard model. We consider a specific model for its generation, by a simple perturbative out of equilibrium decay of Higgs like scalar fields with CP-violating Yukawa couplings to the standard model leptons. We show that, because of the much enhanced expansion rate, such a mechanism can easily produce an adequate asymmetry from scalars with masses in the range $1 - 100\,\mathrm{TeV}$. Kinetic mode domination is strongly favoured because it evades large entropy release which dilutes the asymmetry. We also discuss briefly the effect of the abelian hypercharge anomaly.

Keywords: Baryogenesis, Physics of the Early Universe.

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Contents

1. Introduction

Until a few years ago cosmology with scalar fields was almost synonymous with cosmological inflation. Recently there has been an enormous upsurge in interest in the possibility that scalar fields can play an important role in the dynamics of the Universe at recent epochs, mainly due to the observations of the apparent magnitudes of distant supernovae [1] which may be explained by the presence of such a component [2]. In this context it is certainly interesting to consider what the role of such fields can be at other epochs, and in particular how their behaviour between the end of inflation and their reappearance today might influence cosmology in the intervening period. This question is also related to the 'fine-tuning' problem associated with such scenarios: how is it that such a field can give a significant contribution to the energy density today starting from a natural set of initial conditions after inflation? This apparent problem is in fact resolved in a wide class of potentials [3, 4] which generically have the property that in some part of the potential they may support modes which are dominated by the kinetic energy of the scalar field, so that their energy density scales away faster than that in the radiation, i.e. $\rho_\phi \propto a^{-n}$, with $4 < n \leq 6$, where a is the scale factor. In principle there is no reason why such modes cannot initially dominate over the radiation component, and in certain specific models this is realized. The main important observational constraint is that such domination must terminate by the nucleosynthesis epoch, when the expansion law must be that given by radiation domination with the standard model degrees of freedom. There may be an additional contribution which, conservatively, must be less than about 20% of the total [5].

More generally we can consider the question of the cosmology of the Universe between the end of an inflationary epoch and the entry into radiation domination before nucleosynthesis. For a transition from a scalar field dominated cosmology to occur the energy in the scalar field must either decay (directly or indirectly) into standard model particles — as in standard reheating scenarios [6] — *or* it must redshift away more rapidly than the radiation. Or some combination of the two can occur. In the former case any scaling less rapid than that during inflation ($n > 2$, or equivalently an equation of state $p_\phi = w_\phi \rho_\phi$, with $w_\phi > -1/3$) can be envisaged, with the case $n = 3$ corresponding to the most standard reheating during the oscillation of the inflaton about the minimum of a quadratic potential. There is a continual release of entropy until the radiation dominated epoch, leading to a dilution of most relevant physical quantities sourced during the scalar dominated phase. In the latter case, which corresponds to domination by the kinetic energy of a homogeneous scalar field (or equivalently to an equation of state $p_\phi = w_\phi \rho_\phi$, with $w_\phi > 1/3$) the scalar field simply redshifts away until it becomes the sub-dominant component. There is no entropy release, and correspondingly a coherent energy remains in the scalar field which, given an appropriate potential (the 'self-tuning' potentials of [3], or the 'tracking' potentials of [4]) can become relevant again at late times [7, 8, 9].

In [10, 8] we have considered in a generic way the effect of a change in the expansion rate prior to nucleosynthesis on models of electroweak baryogenesis,[1] in particular on the effect on the sphaleron bound and the 'no-go' theorem for electroweak baryogenesis in the case of a second order phase transition. As concrete realizations of such cosmologies we considered models which go through an epoch after inflation — which, following [10] we termed 'kination' — of domination by a kinetic mode of a scalar field. This occurs most naturally in a model in which the universe 'reheats' not by the decay of the inflaton, but by gravitational particle creation at the end of the de Sitter phase [13, 7]. In a recent paper [14] it has been observed that, for low (sub-electroweak) reheat temperatures in more traditional models of reheating — in which the inflaton decays while oscillating in a mode with matter scaling after inflation — the effects discussed in [10, 8] on electroweak cosmology also result. There is in this case an even larger relative boost to the expansion rate (see below), but a very large entropy release which tends to undo any of the enhancing effects of the

[1] The effects on dark matter freeze-out can be inferred from the work of [11, 12], who studied mainly modifications associated with anisotropy in the expansion.

greater expansion rate. In [15] one of us (TP) has considered the general case of a decaying inflaton evolving in a mode scaling as $1/a^n$, and shown that, while the same larger boost to the expansion rate occurs as in the $n = 3$ case of [14], the entropy release problem is greatly reduced as the kinetic mode $n = 6$ limit is attained.

Here we concentrate on another aspect of such alternative cosmologies, which is a simple consequence of the observation which has been made in [16, 8, 14]: Because of the enhanced expansion rate, the right-handed electrons of the standard model may remain out of equilibrium until a temperature below the electroweak phase transition. It is well known that asymmetry in right-handed electrons — because of their late equilibration time — may be important from at least two points of view:

- since right-handed electrons couple to other particles in the standard model with only an extremely small Yukawa coupling, they remain out of equilibrium in an expanding Universe until relatively late — in the standard radiation dominated cosmology until $T \gtrsim 20\,\mathrm{TeV}$ [17, 18]. A pre-existing baryon asymmetry can survive the effect of standard model anomalous processes - which violate $B + L$ and are unsuppressed until the electroweak phase transition — only if there are non-zero CP odd conserved global charges when they are operative. In the absence of such charges the equilibrium attained will be CP invariant with zero baryon number. As noted in [18] above[2] $20\,\mathrm{TeV}$ right electron number e_R is in fact such an effective charge, and as a result other global charges like $B-L$ can be violated until close to this scale. This leads [18] to a very significant reduction in the bounds on $B - L$ violating interactions in grand unified theories with the structure appropriate for them to preserve baryon asymmetry. Here the consequences are much simpler and more dramatic: If the e_R remain out of equilibrium all the way until the electroweak scale, a baryon number will result from this due to the $B + L$ violating processes. When the electroweak scale is reached this baryon number will simply be frozen when the $B+L$ violating processes abruptly switch off. This will be the case irrespective of whether there is primordial B or L (or $B - L$), and irrespective of whether these charges are violated or conserved. Just like in the case of electroweak baryogenesis all the non-trivial physics required is in principle present in the standard model. The problem of baryogenesis then becomes reduced solely to the generation of the source right handed electrons prior to the electroweak scale at which the $B+L$ violating processes become suppressed. We dub this process 'electrogenesis'. Electrogenesis takes place also in the mechanisms discussed in [17, 18].

- The effective conservation of e_R in the early Universe due to the fact that its perturbative decay channel is out of equilibrium is not exact, because the e_R

[2]The scale quoted in [18] is $10\,\mathrm{GeV}$. The increase by a factor 2 is due to a tighter bound on the Higgs mass.

charge has an axial anomaly under the U(1) of hypercharge. There are no degenerate vacua as in the non-abelian case, but there are finite energy modes of the U(1) field with Chern-Simons number which can 'eat' the charge. In fact, as discussed in [16, 19], this leads to an instability at finite density to the formation of long wavelength modes of hypermagnetic field. When these modes come inside the horizon they can evolve during the time in which the right electron number is without its perturbative decay channel. Here this scenario will be modified as a result of the change in the expansion rate, since the perturbative channel does not come into play until the electroweak scale, at which time a first order phase transition may produce the turbulence needed to amplify the produced seed magnetic fields.

2. Scalar fields and the expansion rate after inflation

The inflationary solutions for scalar fields represent only one part of a much wider range of possible behaviours of the energy density in the zero modes of scalar fields. The full range can be characterized by the equation of state for a (real) scalar field which is determined by the relative weight of the kinetic and potential energy (see [10, 8] for a discussion):

$$p_\phi = w_\phi \rho_\phi \,, \qquad w_\phi = \frac{\frac{1}{2}\dot{\phi}^2 - V(\phi)}{\frac{1}{2}\dot{\phi}^2 + V(\phi)} \,, \qquad \rho_\phi \propto a^{-3(w_\phi+1)} \,. \tag{2.1}$$

The limit of potential energy domination gives inflation, with $\rho_\phi \approx$ const., while the opposite limit of complete kinetic energy domination gives the most rapid possible red-shifting of the energy to be $\rho_\phi \propto 1/a^6$. While inflation is associated with flat potentials (satisfying 'slow-roll' conditions), the latter limit is associated with steep potentials.[3] A particularly useful 'yard-stick' of flatness/steepness is the simple exponential potential

$$V_{\rm exp}(\phi) = M_P^4 e^{-\lambda\phi/M_P} \,, \tag{2.2}$$

where $M_P = 1/\sqrt{8\pi G} \approx 2.4 \times 10^{18}\,$GeV is the reduced Planck mass (and the origin of ϕ has been chosen to give the simple normalization). This potential in fact has an attractor solution [20] for any $\lambda^2 < 6$ in which the energy density scales as $1/a^{\lambda^2}$, and as $1/a^6$ for $\lambda^2 > 6$. A potential with a varying slope, e.g. the inverse power-law potential of [21, 22] $V \sim M_P^{4+\alpha}/\phi^\alpha$ then supports a kinetic mode at small ϕ, but an inflationary type (or 'quintessence') mode at large values of the field. Alternatively an oscillating mode about the minimum of a potential $V \sim \lambda_\alpha \phi^\alpha$ gives a broad range of scalings with $\rho_\phi \propto a^{-6\alpha/(\alpha+2)}$ [23], producing thus the familiar matter scaling when $\alpha = 2$ and radiation scaling when $\alpha = 4$.

[3] The exception is a flat direction with no associated potential energy e.g. a Goldstone direction associated with a broken exact global symmetry, which only has pure kinetic modes.

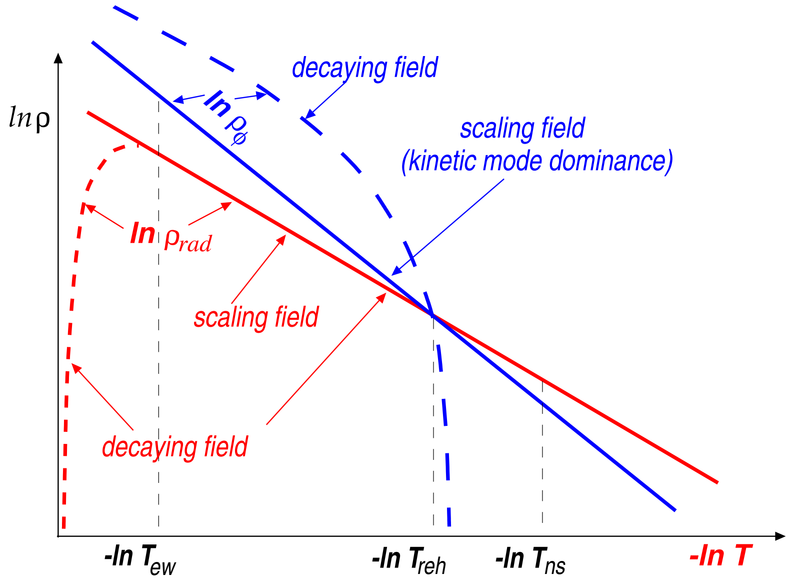

Figure 1: Evolution of energy density in radiation and the dominant scalar field as a function of temperature. Two cases are illustrated: *Model (A)* in which the dominant scalar component scales faster then radiation, but does not decay (*solid lines*), and *Model (B)* in which the scalar field decays (*dashed lines*).

In [8] we discussed several ways in which a period of kinetic mode domination (which, following [10], we termed 'kination') could come about after inflation.[4] We considered only the case in which the relevant field (the 'kinaton') did not decay itself, and discussed two possible sources for the radiation in the Universe: The entropy associated with particle creation during the de Sitter phase (see below), or a more conventional source in the decay of a distinct inflaton field. In the latter case specific conditions need to be satisfied by the 'kinaton' field to allow it to dominate over the energy produced by the inflaton, whereas in the former the kinaton and the inflaton are one field and the domination by the kinetic mode for a period is a built-in and necessary feature.

Our concern in this paper is not the inflationary model building aspect of the problem, but rather the problem of 'electrogenesis' in this kind of cosmology, as well as in the more conventional reheating models discussed in [14, 15]. For the sake of clarity and simplicity we limit ourselves here to two definite and simple models with scalar field dominance continuing until temperatures just above the nucleosynthesis scale, exemplifying these two types of different cases:

Model (A): the inflaton rolls after inflation into a steep potential in which the field rolls in a kinetic mode, so that the energy density scales as $1/a^n$ where $n > 4$ (see figure 1). The field is assumed to be very weakly coupled and the only radiation

[4]Such kinetic modes have also recently been used to propose a solution to the cosmological moduli problem [24].

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present is the very sub-dominant component due to particle creation at the end of the preceeding de Sitter phase. The latter has a characteristic energy density H_I^4, where H_I is the expansion rate at the end of inflation, so that initially $\rho_{\rm rad}/\rho_\phi \sim (H_I/M_P)^2$. Provided the inflaton scales faster than radiation it will become subdominant at a subsequent time and the transition to radiation domination is achieved without any decay of the field [13, 7]. Requiring that this transition occurs before nucleosynthesis gives an absolute lower bound on H_I, which for a pure (or almost) kinetic mode scaling as $1/a^6$ results is $H_I \gtrsim 10^7\,{\rm GeV}$ (see [8]). As noted by Spokoiny [7], for an appropriate potential the field can again dominate in a slowly scaling mode at late times. This kind of model has been dubbed 'quintessential inflation' and studied in more detail in [9, 25, 26]. See also [27] for a discussion of some of the potential problems of this type of model in certain contexts, in particular in supergravity models. Taking the reheat temperature $T_{\rm reh}$ to be defined[5] as that when $\rho_\phi \simeq \rho_{\rm rad}$, it is easy to infer[6] that above this temperature we have

$$H = H_{\rm rad}\left(\frac{T}{T_{\rm reh}}\right)^{\frac{n-4}{2}}\,, \tag{2.3}$$

where $H_{\rm rad} \simeq 1.4\times 10^{-16}(T^2/100\,{\rm GeV})$ is the standard radiation dominated evolution of the expansion rate, and n gives the scaling of the energy density in the dominant scalar mode $\rho_\phi \propto 1/a^n$, with clearly the largest enhancement of the expansion rate for the limit $n = 6$. The constraint that the energy density in the scalar field be less than about 20% at nucleosynthesis requires that $T_{\rm reh} \gtrsim 5^{1/(n-4)}T_{\rm ns}$ (and $T_{\rm ns} = 1\,{\rm MeV}$). Here we are interested in the case when right electrons are out of equilibrium at the electroweak scale, which corresponds therefore to the upper bound

$$T_{\rm reh} < T_{\rm ew}\left(\frac{H_{\rm ew}}{\Gamma_{e_R}}\right)^{\frac{2}{n-4}}\,, \tag{2.4}$$

which for the optimum case ($n = 6$) becomes

$$T_{\rm reh} < T_{\rm ew}\frac{H_{\rm ew}}{\Gamma_{e_R}} \sim \frac{T_{\rm ew}}{200} \sim 0.5\,{\rm GeV}\,, \tag{2.5}$$

where we have made use of the fact that the interaction rate for right electrons through their Yukawa coupling is $\Gamma_{e_R} \approx 10^{-13}x_{e_R}^2 T$ [18], and we took $x_{e_R} \equiv m_H/2T_{\rm ew} \sim$

[5] Note that in these models the Universe is strictly speaking not 'reheated' at all — the entropy is left behind at the end of the de Sitter phase and the important process is the red-shifting away of the dominant energy in the inflaton. Here we adopt the standard definition of 'reheat temperature' as used in standard reheating models. In [8] we used 'reheat temperature' to mean the temperature of the radiation when it first thermalizes, which is far higher ($\sim 0.1H_I$) than the 'reheat' temperature as defined here. What we now call $T_{\rm reh}$ is denoted $T_{\rm k,end}$ ('end of kination') in [8].

[6] For simplicity we neglect here and elsewhere the small reheating factors associated with particle decouplings.

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0.5 correspondinding to the current lower bound on the Higgs mass m_H. In terms of the expansion rate the bound (2.5) corresponds to the requirement of a boost by about 200 times in the expansion rate at the electroweak scale relative to that in the standard radiation dominated cosmology.

Model (B): the inflaton evolves after inflation into a potential, in which it rolls or oscillates, scaling as $1/a^n$ with $6 \geq n \geq 3$. The dominant source of entropy comes from the decay of the inflaton, which is however sufficiently weakly coupled that reheating occurs between the electroweak scale and the nucleosynthesis scale. The energy density-temperature dependence for this case is illustrated in figure 1. The phase we are discussing corresponds to the 'preheating' phase [28] of inflationary models with the usual mechanism of reheating from inflaton decay, either in an oscillatory mode (with $n = 3$ for a ϕ^2 potential) or a rolling mode. We assume here for simplicity perturbative reheating, but note that the nonperturbative decay channels of narrow resonance may also be considered. A particular realization of the case of a rolling mode is given by the 'NO' models of [27], with reheating by a mechanism discussed in [29] ('instant preheating').

It is quite easy to show [15] that in *Case (B)* the expansion rate as a function of the temperature is independent of the equation of state (2.1), i.e. the following *universality in scaling* in the expansion rate holds

$$H = \frac{5-3w_\phi}{6}\frac{\rho_r}{\Gamma_\phi M_P} = H_{\rm rad}\left(\frac{T}{T_{\rm reh}}\right)^2 , \tag{2.6}$$

which implies that, for a reheat temperature below the electroweak transition, the expansion rate is enhanced by $(T/T_{\rm reh})^2$ with respect to the standard rate $H_{\rm ew} \equiv H_{\rm rad}(T_{\rm ew})$. The condition that the right electrons remain out of equilibrium until the electroweak scale is in this case

$$T_{\rm reh} < T_{\rm ew}\left(\frac{H_{\rm ew}}{\Gamma_{e_R}}\right)^{1/2} \sim \frac{T_{\rm ew}}{15} \sim 5\,{\rm GeV}\,, \tag{2.7}$$

which again corresponds to the same minimal boost in the expansion rate by a factor of about 200. The extra increase in the expansion rate as a function of temperature compared to the first case is due to the 'leaking' of the scalar field energy into the radiation. Note that the lower bound on $T_{\rm reh}$ is in this case $T_{\rm reh} \gtrsim 2\,{\rm MeV}$.

Our interest here finally is in the ratio of baryon number to entropy, and so will need to include the dilution effect of this entropy production subsequent to the scale $T_{\rm dec}$ at which the baryon number, or in fact the source for it, right electron number is produced. As discussed in [15] the entropy per comoving volume $S_{\rm com}$ scales as $a^3T^3 \propto T^{-3(8-n)/n}$ since $a \propto t^{2/n} \propto T^{-8/n}$. Thus the dilution factor $f_{\rm dil}$ due to entropy production between the two scales is

$$f_{\rm dil} \simeq \left(\frac{T_{\rm dec}}{T_{\rm reh}}\right)^{\frac{3(8-n)}{n}} . \tag{2.8}$$

Thus there is a very significant difference between the case of the matter scaling (considered in [14]) giving $f_{\rm dil} \simeq (T_{\rm dec}/T_{\rm reh})^5$ and that of the kinetic mode limit with $f_{\rm dil} \simeq (T_{\rm dec}/T_{\rm reh})$. The origin of this difference can be easily understood: a scalar kinetic mode gets rid of most of its energy by the rapid red-shifting.

We now turn to the effect of these modifications to the pre-nucleosynthesis expansion rate on the generation of a baryon asymmetry from an e_R asymmetry.

3. From e_R to a baryon asymmetry

Before discussing the generation of right electron asymmetry in these cosmologies, we discuss the conversion of such an asymmetry to a baryon asymmetry when $B+L$ violating processes are active (adapting the discussion of [18] to the present case). 'Conversion' is in fact a little misleading as these processes of course only act on the left-handed fermions: As will become more explicit now the physics of the creation of the baryon asymmetry is that the right electrons carry the gauge charge hypercharge, which is globally zero and exactly conserved. When there is net hypercharge in the e_R sector, there must be also a compensating hypercharge in the rest of the particles. When this is non-zero the $B+L$ violating processes minimize the free energy with a non-zero baryon number.

We follow a standard procedure and consider the equilibrium abundance of baryon number subject to the constraints imposed by the charges conserved by the fast interactions (which are in equilibrium at that time). Because baryon number violation freezes out at the electroweak scale, when the sphaleron processes become suppressed, this is the scale at which we need to calculate baryon number. Above the electroweak scale the rate of the B-violating processes is mediated by the symmetric phase sphaleron transitions, which are unsuppressed: $\Gamma_{\rm sph} \approx 25\alpha_w^5 T \sim 10^{-6}T$ [30], so that they will have time to equilibrate above the electroweak scale in the models we are discussing (cf. eqs. (2.3) and (2.6)). In fact we shall assume for simplicity that the expansion rate is such that the e_R are the only standard model degrees of freedom out of equilibrium. Within the scenarios we are discussing this is not necessarily the case, as the rate could be enhanced in principle enough to take also other heavier particles out of equilibrium. For example the μ_R has a Yukawa coupling larger by about 10^2 and therefore a decay rate faster by a factor of 10^4, while with a reheat temperature sufficiently close to the nucleosynthesis scale the expansion rate may be boosted in *Model A* by almost as much as $T_{\rm ew}/T_{\rm ns} \sim 10^5$, and in *Model B* by $(T_{\rm ew}/T_{\rm ns})^2 \sim 10^{10}$ times (with of course the correspondingly large entropy release factor).

With the following hierarchy of couplings

$$\Gamma_{e_R} \ll H \ll \Gamma_{\rm sph}, \Gamma_{\mu_R}, .. \tag{3.1}$$

the appropriate equilibrium calculation of baryon number is particularly simple. In the standard model the only conserved charges are the gauge charges, e_R and $\frac{1}{3}B - L_i$, where the latter is the baryon minus lepton number in each generation. We will make the slight simplification of assuming only total $B-L$ as conserved (which would be appropriate at this scale in certain models including neutrino mass matrices), which leads to minor numerical changes to the results quoted here. (We refer the reader to [8] where the full set of constraint equations can be found.)

To arrive at the set of constraint equations one expresses the charge densities in terms of particle densities n_α using $n_\alpha - \bar{n}_\alpha = (T^2/6)k_\alpha\mu_\alpha$, where[7] $k_\alpha = 1(2)$ for fermions (bosons) and μ_α is the chemical potential for a species α. Further μ_α can be re-expressed in terms of the chemical potentials for charges Q_A as follows: $\mu_\alpha = \sum_A q_\alpha^A \mu_A$, where q_α^A is the A-charge of the α species. With this we can have (cf. [8]) the following constraint equations while the baryon number is out of equilibrium:

$$\begin{aligned} Y &= \frac{T^2}{6}\left[(10+n)\mu_Y + 8\mu_{B-L} + \mu_{e_R}\right] \\ B-L &= \frac{T^2}{6}\left[8\mu_Y + 13\mu_{B-L} - \mu_{e_R}\right] \\ e_R &= \frac{T^2}{6}\left[-\mu_Y - \mu_{B-L} + \mu_{e_R}\right]. \end{aligned} \tag{3.2}$$

Here we used the hypercharge assignments such that $Q = Y + T^3$, where Q denotes the electric charge and T^3 the isospin. We have not written the second linearly independent gauge charge explicitly, as choosing it as T^3 it is simply proportional to its own chemical potential, and so trivially drops out of the equations when we impose $T^3 = 0$. The baryon number B can itself be expressed in terms of the relevant chemical potentials as

$$B = \frac{T^2}{6}\left[2\mu_Y + 4\mu_{B-L}\right]. \tag{3.3}$$

The gauge charge Y must be zero, and then given any value of the global conserved charges eqs. (3.2) can be solved to give the baryon number (3.3). When $B-L$ is conserved we thus have

$$B = \frac{2(9+2n)}{59+12n}\, e_R + \frac{2(11+2n)}{59+12n}\,(B-L) \tag{3.4}$$

and see that e_R is an almost equally strong source for baryon number as is $B-L$. Indeed, as n changes from $n=1$ to $n=\infty$, the coefficient of e_R changes from 0.31 to 1/3, while that of $B-L$ from 0.32 to 1/3. Hence eq. (3.4) may be quite well approximated by

$$B \approx \frac{1}{3}\left(e_R + (B-L)\right). \tag{3.5}$$

[7] We use here the massless approximation, to which there will be small corrections due to thermal masses. Note that we also assume the right electron distributions can be described by a chemical potential, which is justified given their relatively fast elastic scattering rate through weak hypercharge processes $\sim 10^{-2}T$ [31]

Thus if $B - L$ is conserved by all interactions after inflation, it is zero and remains zero, but the final baryon number, in contrast to the usual radiation dominated universe, is now non-zero and simply proportional to the original e_R asymmetry. Indeed therefore we see explicitly that no B violation other than that of the anomalous processes of the standard model is required to produce it.

While the latter is the case which will interest us, it is interesting to note that the result that one obtains a non-zero baryon number from e_R is very robust, and is relatively insensitive to whether the other charges are violated. Indeed it's easy to see that if e_R is the *only* conserved charge — for example if $B - L$ is violated by some interactions all the way down to the electroweak scale — the net baryon number is still non-zero. Solving the reduced set of constraint equations for Y and e_R only, with μ_{B-L} set to zero, we find

$$B = -\frac{2}{11+n}\, e_R\,, \tag{3.6}$$

which is slightly smaller and of the opposite sign than the result in eq. (3.4).

We conclude that, irrespective of constraints on the value of $B - L$ and assumption on whether $B-L$ is conserved, a right-handed electron asymmetry is reprocessed into a baryon asymmetry of the same order.

4. Electrogenesis: creation of the e_R asymmetry

We now consider explicitly the production of a right-handed electron asymmetry prior to the electroweak scale. In the standard radiation dominated cosmology right electrons have been understood to be of interest because of their capacity to protect a baryon asymmetry from erasure [17, 18]. Thus their generation has been considered in the context of theories which also produce such a primordial (total) baryon or lepton asymmetry, and thus typically the scale characterizing their generation is very high, around the GUT scale or in the case of leptogenesis as low as 10^{10} GeV [32]. In the present context right electrons are in their own right adequate sources for baryogenesis by reprocessing with standard model $B + L$ violation. Given that the physics required to generate them is CP violating only, and thus potentially can be associated in simple ways with much lower energy scales, it is certainly of interest to consider mechanisms which can produce them quite independently of (total) B or L violation beyond the standard model.

In fact in the cosmologies being considered one is forced to seek such different mechanisms of e_R generation for another very simple reason which we have not drawn attention to so far: The maximum temperature $T_{\max}$ attainable after inflation in these cosmologies is in fact much lower than in the standard radiation dominated cosmology. Given the requirements of $T_{\rm reh}$ in (2.4) and (2.7), we can bound the temperature above by extrapolating the expansion rate to the point $H \sim 0.1T$. For model A this corresponds to $T_{\max} \lesssim 10^8\,\mathrm{GeV}$, while for model B it gives $T_{\max} \lesssim 10^6\,\mathrm{GeV}$. Above this point thermodynamic temperature can have no meaning as the age of the Universe is shorter than the equilibration time of any process. Thus any mechanism which in the ordinary radiation dominated scenario relies on temperatures being reached higher than this is not applicable, and we must seek mechanisms which operate at a lower temperature.

Here our aim is not to be exhaustive about possible mechanisms, but rather to study an explicit model which produces an e_R asymmetry sufficiently large to source the observed baryon asymmetry, in these cosmologies. Given that in principle all the elements are present in the standard model itself, it is natural to ask — just as one does in the context of electroweak baryogenesis - whether it alone might suffice. While in the standard radiation dominated cosmology the standard model has apparently insurmountable problems on two fronts [33] — the sphaleron bound and the inadequacy of standard model CP violation — here the former does not provide a significant constraint. All we require is that the right electron come into equilibrium after the $B+L$ violation goes out of equilbrium, not that the sphaleron rate exactly at the transition be smaller than the expansion rate. So could the standard model with it CP violation produce the e_R asymmetry? Given that its production can only come about through the same Yukawa coupling channel, the answer would seem to be definitively in the negative. In general however the question in these cosmologies can be framed more generally given that the expansion rate can change enormously: Is it possible to generate some CP odd charge (not necessarily e_R) which is conserved on a time scale longer than that associated with the $B+L$ violating processes in the unbroken phase? We will return briefly to this question in the conclusion.

Here, just as one does in the context of baryogenesis models, we add some extra CP violating physics in the scalar sector. We study a simple out of equilibrium decay of scalar particles with CP violating decays. Interestingly we find that, again because of the modified expansion rate prior to nucleosynthesis, the mass of these scalars need not be so far above the electroweak scale for the mechanism to work. This suggests that the kind of mechanism for 'electrogenesis' we discuss may be implementated successfully in other theories with additional scalars particles at scales not far above this scale, with testable signatures at accelerators. We will return to this point in our conclusions.

4.1 The model

The additional particle content we assume over the standard model (and the inflaton) is a set of Higgs-like scalar doublets Φ^a coupled to the standard model leptons through a Yukawa type interaction i.e. with interaction lagrangian

$$\mathcal{L}_{\mathrm{add}} = -h^a_{ij}\Phi^a\bar{\psi}_{iL}\psi_{jR} + \mathrm{h.c.} \tag{4.1}$$

where the couplings h^a_{ij} are CP violating, i.e. $\mathbf{h^a}^\dagger \neq \mathbf{h^a}$, where $\mathbf{h^a}$ is the matrix of couplings. While in principle CP violation does not mandate a matrix of couplings, but only a coupling to the right electron itself with a complex phase unremovable by phase transformations on the whole lagrangian, we will require the flavour mixing structure and the existence of at least two such scalars in order to implement the generation of a CP violating asymmetry. The strongest constraints on the masses and the couplings of such scalars come from the fact that they are flavour changing. For leptons the strongest constraint of this type comes from the bounds on the decay $\mu \to e\gamma$ [34, 35]. For couplings h of order one this requires masses $M_\Phi \gtrsim 100\,\mathrm{TeV}$, with the branching ratio for this process going parametrically as $h^2_{\mu\tau} h^2_{e\tau}(M_W/M_\Phi)^4$ so that much smaller masses can be permitted if the couplings have the a hierarchy like that in the standard model Yukawa couplings [35, 36].

4.2 The out of equilibrium conditions

We consider here a simple out of equilibrium decay scenario for these particles, very analogous to that which occurs in standard GUT scale baryogenesis scenarios [6]. It is possible that nonperturbative decay mechanisms may be operative and work just as well, but we limit our treatment here to the simpler perturbative case. The perturbative decay rate for Φ can be well approximated by

$$\Gamma_{\phi,\,\mathrm{pert}} = \frac{|\mathbf{h}|^2}{8\pi} E_\phi, \tag{4.2}$$

where E_ϕ is the energy of Φ, $|\mathbf{h}|^2 = \mathrm{Tr}(\mathbf{h}\mathbf{h}^\dagger)$ and we have assumed the energy of the Φ is much greater than that of the produced fermions (e.g. in the case $m_i = m_j = m$ there is a simple suppression $E_\phi \to [E^2_\phi - 4m^2_\psi]^{1/2}$).

Before considering the production of a CP asymmetry we first discuss the out of equilibrium condition. When the particles decay, with rate given by (4.2), the reverse process (or any other one) creating them must be suppressed. This is fulfilled if the temperature of the plasma at the time of decay is well below the mass scale of the scalars i.e.

$$M_\Phi > T \qquad \text{when} \qquad \Gamma_\phi \sim H\,. \tag{4.3}$$

Equations (2.3) and (2.6) give the boost to the expansion rate with respect to the radiation dominated case as $(T/T_{\mathrm{reh}})^p$, where $p = 1$ for kinetic mode domination($n = 6$), and $p = 2$ for a dominant decaying component. Making use of this and eq. (4.2), we infer that the constraint (4.3) can be re-expressed as

$$M_\Phi > T_{\mathrm{dec}} > (70 g_*)^{-1/2(1+p)} \left[|\mathbf{h}|^2 M_P T^p_{\mathrm{reh}}\right]^{\frac{1}{1+p}} \qquad (0 \leq p \leq 2)\,, \tag{4.4}$$

where T_{dec} is the temperature at which Φ decays, and we have used $H_{\mathrm{rad}} = (\frac{\pi^2 g_*}{90})^{1/2} \times T^2/M_P$ (where $M_P \approx 2.4 \times 10^{18}\,\mathrm{GeV}$). For the case of radiation domination ($p = 0$) this gives $M_\Phi > T_{\mathrm{dec}} > 10^{16}|\mathbf{h}|^2\,\mathrm{GeV}$, where we took $g_* \sim 10^3$. Given that in these scenarios the asymmetry is generated by, at the very least, the interference between a tree-level and one-loop diagram, it is always suppressed by some small numbers times at least a square of the couplings $\mathbf{h}$, and typically by higher powers of the couplings. To produce a significant asymmetry one cannot have the coupling too small, and conversely one needs the scalar field to have a mass not so far below the GUT scale.

For the cosmologies we are considering these bounds change very considerably. In model A in which the Universe is dominated by a kinetic mode ($p = 1$) the constraint (4.4) relaxes to

$$M_\Phi > T_{\rm dec} > 3\times 10^6\,{\rm GeV}|\mathbf{h}|\left(\frac{T_{\rm reh}}{T_{\rm ns}}\right)^{1/2} > 5\,|\mathbf{h}|\times 10^6\,{\rm GeV}. \tag{4.5}$$

where we took $T_{\rm ns} = 1\,{\rm MeV}$ and $g_* \sim 10^3$. This should be compared with the energy scale H_I which characterizes this model at the beginning of the post-inflationary epoch. For a reheat temperature $T_{\rm ns}$ and a pure $n = 6$ scaling after inflation one finds [8]

$$H_I \sim 10^7\,{\rm GeV}\left(\frac{{\rm T}_{\rm reh}}{{\rm T}_{\rm ns}}\right)^{1/2} \tag{4.6}$$

Thus the M_Φ can be sufficiently light that they are produced by gravitational coupling in this mechanism along with all the other lighter ($m < H$) degrees of freedom. A little subsequently, at a temperature $T \sim 0.1H_I$ the strongly interacting degrees of freedom can begin to equilibrate (and define a real thermodynamic temperature), while the M_ϕ can decay without ever coming into equilibrium. For smaller values of the coupling h ($< 10^{-2}$) there may be some time for weak force mediated annihilation processes (with rate $\sim \alpha_w^2 T$) to act, and in this case the initial e_R number density at the time of decay will be reduced somewhat relative to their initial value.

In the case of Model B, when the dominant component decays, we have $p = 2$ so that eq. (4.4) gives an even milder bound on the mass of Φ:

$$M_\Phi > T_{\rm dec} > 2\,|\mathbf{h}|^{2/3}\left(\frac{T_{\rm reh}}{T_{\rm ns}}\right)^{2/3}\,{\rm TeV} > 3\,|\mathbf{h}|^{2/3}\,{\rm TeV}. \tag{4.7}$$

In this case therefore the out of equilibrium condition may in some cases (for sufficiently low $T_{\rm reh}$) provide an even weaker constraint on their masses than accelerator constraints from the flavour changing processes they can mediate. More generally, it is certainly interesting to note that the mass scale is sufficiently low that models may be viable in which the scalars are the scalar partners of the standard model particles. We will return to this point in our conclusions. In models of type B therefore we envisage that the scalars are created during reheating in similar quantities to the other degrees of freedom (i.e. equipartition of energy), and that they subsequently decay when the temperature $T_{\rm dec}$ is reached. As in case A one would need to consider carefully the different cases (depending on $|\mathbf{h}|$) in which the weak interactions can or cannot play a role in reducing the particle anti-particle asymmetry in Φ before this decay occurs. One feature of (4.7) should immediately be noted, however, and we will return to it below: The entropy release of these models which is of relevance in the present case is that which occurs between the time of production of the e_R asymmetry, $T_{\rm dec}$, and $T_{\rm reh}$. From (4.7) it follows that

$$\frac{T_{\rm dec}}{T_{\rm reh}} > 2\,|\mathbf{h}|^{2/3}\left(\frac{\rm TeV}{T_{\rm ns}^{2/3}T_{\rm reh}^{1/3}}\right) \gtrsim 10^5\,|\mathbf{h}|^{2/3} \tag{4.8}$$

where the latter equality follows from (2.7). When it comes to producing a final baryon asymmetry this constraint will make it very difficult for models with any scaling much faster than the kinetic mode to produce a reasonable final baryon to entropy ratio. We will return to this below.

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4.3 Generation of the asymmetry

We now turn to the production of the asymmetry through the out of equilibrium decay of these scalar fields.

As in any out of equilibrium decay scenario, one must go beyond the tree-level decay and consider the interference between tree-level and higher order diagrams, to produce any net CP violating effect. Further since CPT violation implies the equality between the total decay rate for particles and anti-particles, we need at least two channels containing different e_R number in order to be able to produce the asymmetry. It is to this end that we have taken the Φ^a scalars to couple to more than one generation. Further to produce a CP violating effect at one loop level it is not difficult to see that we need at least two scalars, just as one requires two heavy bosons in simple GUT scenarios (cf. [6]). Here we do not try to be exhaustive in our consideration of the diagrams which give dominant contributions in different parts of parameter space. In figure 2 are shown the two diagrams we consider here, for the decay channel $\Phi \to \bar{e}_R \Psi^i_R \Phi'$ where the Φ' is assumed to be the lighter of the (at

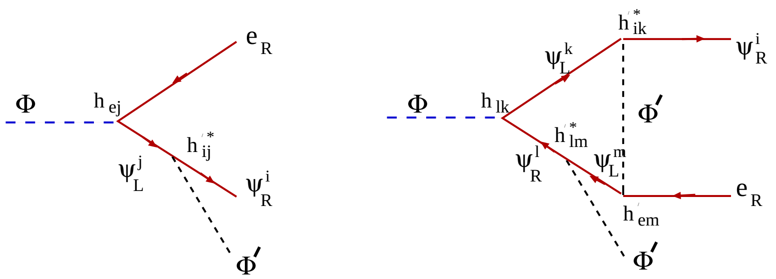

Figure 2: Tree and one loop diagrams for the three body decay $\Phi \to \bar{e}_R \Psi^i_R \Phi'$, with the appropriate couplings at the vertices. We assume that Φ is heavier than Φ'. When the second outgoing lepton is a μ or τ lepton the process produces net e_R number.

least) two scalars. Provided the second fermion Ψ_R^i is of one of the two heavier lepton flavours, the process violates right electron number. The rate of the corresponding anti-particle decay does not cancel if the CP violating interference terms between the two diagrams have a pure complex part. Summing over the internal fermions we have (see [37] for somewhat similar cases) that the net CP violating effect creating net e_R number compared to the tree-level decay is

$$\frac{\Gamma_{\Phi\to\bar{e}_R\Psi_R^i\Phi'} - \Gamma_{\bar{\Phi}\to e_R\bar{\Psi}_R^i\bar{\Phi}'}}{\Gamma_\Phi^{\rm tot} + \bar{\Gamma}_{\bar{\Phi}}^{\rm tot}} = \epsilon_p \frac{\mathrm{Im}\left[(\mathbf{h}\mathbf{h}'^\dagger)_{ei}(\mathbf{h}'\mathbf{h}^\dagger\mathbf{h}'\mathbf{h}'^\dagger)_{ie}\right]}{\mathrm{Tr}\left[\mathbf{h}\mathbf{h}^\dagger\right]}, \tag{4.9}$$

where $\epsilon_p \sim 10^{-2}$ is the phase space factor (and i is summed over the non-electron indices). There is also another pair of diagrams which differ only in that the Φ' emission on the external e_R leg, which gives the (4.9) with the indices interchanged. In the result we see clearly that for an effect at this order we indeed need two scalars as when $\mathbf{h} = \mathbf{h}'$ the result vanishes. We note that a diagram in which the Φ' on the external leg is the standard model Higgs could dominate if the $\mathbf{h}'$ couplings are all smaller than the Yukawa coupling of the τ lepton ($y_\tau \approx 10^{-2}$).

We thus write the resultant e_R asymmetry as

$$\frac{e_R}{s}(T_{\rm dec}) \sim \frac{10^{-2}}{g_*} |\mathbf{h}|^4 \, \delta_{\rm CP} \,, \tag{4.10}$$

where $\delta_{\rm CP}$ is proportional to the imaginary part in (4.9). If all the couplings h are of the same order this can be of order one, while if there is a strong hierarchy it will be correspondingly smaller.

4.4 The baryon asymmetry

To arrive at the final baryon asymmetry the cases of model A and model B are quite different. In the former case there is no entropy dilution and the final baryon to entropy ratio is quite simply given by (3.5), i.e. one third of the right electron to entropy ratio. We thus have $n_B/s \sim 10^{-4}|\mathbf{h}|^4\delta_{\rm CP}$, which for roughly equal on and off diagonal couplings in $\mathbf{h}$ and $\mathbf{h}'$ will give a baryon asymmetry of the required size $n_B/s \sim 10^{-10}$ for couplings of order $\sim 10^{-1}$–10^{-2}. Note that this corresponds (from (4.5)) to scalar masses as low as about $M_\phi \sim 100\,\mathrm{TeV}$.

In model B there is the all important entropy dilution factor, so that for the final baryon asymmetry we find

$$\frac{n_B}{s} \simeq \frac{10^{-2}}{g_*} |\mathbf{h}|^4 \, \delta_{\rm CP} \left(\frac{T_{\rm reh}}{T_{\rm dec}}\right)^{\frac{3(8-n)}{n}} \tag{4.11}$$

If we take the constraint given in (4.8) for couplings and $\delta_{\rm CP}$ of order one, we can have a baryon asymmetry compatible with that required for nucleosynthesis only for the case $n \approx 6$, i.e. for the case that the inflaton rolls while decaying in a kinetic energy dominated mode. In this case to satisfy the out of equilibrium condition the mass M_ϕ can be as small as a few TeV. However the constraints on the flavour changing neutral currents need to be carefully considered, but can still be satisfied (e.g. if one of the couplings $h_{e\mu}$ or $h_{e\tau}$ is much smaller than the others). This provides in principle an interesting probe at accelerators of these models in a parameter range which is of interest.

For the standard matter scaling ($n = 3$) [14], or indeed radiation scaling ($n = 4$) during reheating, the entropy dilution factor in (4.11) is much too large to allow the generation of the required baryon asymmetry, and the mechanism we have discussed is not a viable one for baryogenesis in these cases. In a different model it may be possible to relax the condition (4.8) and reduce the dilution factor. This would be appropriate for example if the e_R continues to be created all the way down to the electroweak scale, for example in a scenario in which the Φ particles are themselves directly produced out of equilibrium by the inflaton decay all the way to that scale. Tuning the value of $T_{\rm reh}$ to be just enough to keep the e_R out of equilibrium till that time i.e to satisfy (2.7), the case $n = 3$ gives a dilution by a factor $\sim 10^6$, and the case $n = 3$ by $\sim 10^4$. A fairly copious initial e_R asymmetry must therefore be produced very close to the electroweak scale in order to give the required baryon asymmetry.

The model we have presented here is completely perturbative. It is likely that there are models of non-perturbative decay of a condensate of the Φ field in which the constraints inferred on the Yukawa coupling $\mathbf{h}$ may be relaxed. One simple possibility would be a variant of the well-known Affleck-Dine mechanism [38], with a scalar field Φ charged under right-handed electron number, which oscillates and decays. This may occur for example in supersymmetric extensions of the Standard Model when an $A\Phi^3+$h.c. term is present in the potential for a weakly coupled scalar field. When the field decays it creates a net right-handed electron number which is not suppressed by any coupling constant. The suppression may also potentially evaded without the scalar field being required to carry a right electron charge. This would occur if there were a resonant decay into the electrons, which may in fact occur through precisely the same Yukawa coupling discussed here. Because they do not have the Yukawa coupling suppression, these mechanisms for producing e_R would leave more space for entropy release in models like our model B. However as we have pointed out this will only work in the special case that the e_R is created very close to the electroweak scale, and the temperature at which radiation domination begins $T_{\rm reh}$ lies just far enough (but not too far) below that scale to make keep the e_R out of equilibrium until the electroweak scale.

5. Anomaly and magnetic fields

So far we have neglected the effects related to the abelian anomaly discussed in the introduction. As described in [16] the effect of this term on a finite chemical potential

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μ_R for right electrons is to cause an instability in modes with $k < \mu$. Naively this instability can begin to grow when the corresponding mode enters the horizon (i.e. $\mu \sim H$) , but when the slowing effect of conductivity σ in the plasma is taken into account the criterion becomes $\mu^2/(\sigma) \sim H$. In the standard case, in which the mode must be able to start to evolve before the perturbative e_R decay comes back into play, this corresponds to the effect being important if $\mu/T \gtrsim 10^{-6}$. Here since the expansion rate is such that the e_R come back into equilibrium below the electroweak scale, the lower bound (which depends now simply on the expansion rate at the electroweak scale) will in fact be the same or larger. Moreover, while in the standard case the effect of the perturbative channel is important, and leads to the requirement that a significantly larger asymmetry than the critical one be present initially in order for the source hypermagnetic field to survive to the electroweak scale, here the instability will simply develop all the way to that scale and there will be no damping of the fields due to the appearance of the perturbative channel.

The conclusions we can draw are as follows:

Model A: if the initial e_R asymmetry is less than the critical value for hypermagnetic field generation, all our analysis given above holds. A baryon asymmetry is produced of the same order, and thus for compatibility with the observed asymmetry one must have the corresponding initial value $\mu_R \sim 10^{-10}T$. For initial μ_R larger than the critical value, hypermagnetic field will be generated. The baryon asymmetry generated will depend on the attained value of the chemical potential μ_R (which reduces as the chern simons number grows in the condensate). However, the latter will always be bounded below by the same critical value, and thus the baryon asymmetry also (See also [19] for a discussion of the full evolution of the dynamical equations). Thus we conclude that magnetic field generation from μ_R cannot be attained in this model with an altered expansion rate, since it will always be associated with a baryon asymmetry which is too large. On the other hand, the baryogenesis mechanism from right electrons works perfectly well, and is unaffected by the anomalous effects as the chemical potential is so small.

Model B: the effects of the interplay of the baryon generation and the instability causing the growth of magnetic field are much more difficult to evaluate in this case, and will be model dependent. The fact that the entropy dilution effect mandates a larger initial electron asymmetry, which would then be subject to the instability, suggests that it may be possible to find a model in which both magnetic field and the observed baryon asymmetry could be produced from the right electrons. As was discussed above, the production of the baryon asymmetry would require that a large μ_R chemical potential be produced very close to the electroweak scale. The corollary is that if it produced close to the electroweak scale there will be little time for the instability to evolve and create significant seed fields. One way of getting around this would be in a model with a continuous sourcing of the e_R asymmetry, so that the contribution at earlier time may grow into a magnetic field. The driving chemical potential is, however, itself being constantly diluted by entropy production, and the energy in the resultant field also relative to the background energy density. To see whether seed fields of significant magnitude can survive to the electroweak scale would require detailed study in a model of e_R generation quite different to that we have discussed.

6. Conclusions

We have considered here one aspect of cosmologies in which a scalar field dominates the expansion rate prior to nucleosynthesis. Right-handed electrons may remain out of equilibrium until the electroweak scale, so that if they are generated the $B + L$ processes of the standard model will lead to a non-zero equilibrium density of baryons directly proportional to that in the e_R. We have discussed two kinds of post-inflationary cosmologies in which such a period of scalar field domination can occur : in the first the inflaton rolls away after inflation into a kinetic mode in a steep potential, and the Universe is 'reheated' by the gravitational particle production at the end of the de Sitter phase, while in the second the inflaton rolls into a mode which can have a range of scalings and reheats the Universe itself by decaying sufficiently slowly to give a very low reheat temperature. We studied a specific model for the generation of the right handed electron asymmetry in which there are a set of scalars with CP violating (and flavour changing) couplings to the leptons. We showed that in both scenarios such scalars can decay out of equilibrium at quite low temperatures and produce the desired asymmetry. While our models strongly favoured the case of kinetic mode domination, which have little or no entropy release, we note that in certain very special circumstances which may be satisfied in other models the generation of the observed baryon asymmetry may still be possible in the standard reheating scenario (with matter scaling during the reheating epoch). Finally we considered briefly the effect of the abelian anomaly which destabilizes such charges, and concluded that in the models with kinetic mode domination this effect is unimportant for the baryon number generation, while in the case with large entropy dilution it may be important and might allow the generation of magnetic field as in the case of standard radiation domination.

Finally we return to the question of how this kind of mechanism might be implemented in other particle physics models, in particular in the standard model or popular extensions of it like the minimal supersymmetric standard model (MSSM). In general one need not consider necessarily the generation of right-handed electron number, but the generation of any CP odd charge which is effectively conserved after its creation on a timescale which is longer than the expansion rate of the Universe at the time the sphaleron processes freeze out (at, or shortly after, the electroweak

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phase transition). Given that the expansion rate at the electroweak scale can be enhanced in these models by many orders of magnitude — up to a rate $\sim 10^{-11}T_{\rm ew}$ in models of type A, and $\sim 10^{-6}T_{\rm ew}$ in type B — scenarios can be considered in which many of the lighter degrees of freedom will drop out of equilibrium (for example the lighter right-handed quarks). The problem of baryogenesis then becomes the problem of the generation prior to, or close to, the electroweak scale of such charges using the CP violation available in these theories, within the constraints imposed by these modified pre-electroweak cosmologies.

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