Gravitational leptogenesis from metric perturbations

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In this work we make the observation that the gravitational leptogenesis mechanism can be implemented without invoking new axial couplings in the inflaton sector. We show that in the perturbed Robertson-Walker background emerging after inflation, the spacetime metric itself breaks parity symmetry and generates a non-vanishing Pontryagin density which can produce a matter-antimatter asymmetry. We analyze the produced asymmetry in different inflationary and reheating scenarios. We show that the generated asymmetry can be locally comparable to observations in certain cases, although the size of the matter-antimatter regions is typically much smaller than the present Hubble radius.

I. INTRODUCTION

The excess of matter over antimatter in the universe is one of the longstanding problems in cosmology [1]. This matter-antimatter asymmetry is usually quantified through the ratio of the net baryon number density with respect to the total entropy density whose value measured by the Planck collaboration is $n_B/s = 8.718 \pm 0.004 \times 10^{-11}$ [2]. One of the most interesting proposals for the generation of the baryon asymmetry is leptogenesis [3]. The original implementation of this mechanism relied on the introduction of right-handed Majorana neutrinos in the Standard Model, whose mass term breaks lepton symmetry. This lepton asymmetry is later on converted into baryon asymmetry through electroweak sphaleron processes [4]. For right-handed neutrinos in thermal equilibrium, this mechanism requires a reheating temperature $T_{RH}$ above the right-handed neutrino mass $m_R$ which should satisfy $m_R \gtrsim 10^9$ GeV [5, 6]. For non-thermally produced neutrinos these constraints could be relaxed [7–9].

In [10], an alternative mechanism for leptogenesis was proposed which is not based on the introduction of heavy Majorana leptons. In this gravitational leptogenesis mechanism, lepton asymmetry is generated from the chiral gravitational lepton anomaly already operating in the Standard Model with only left-handed neutrinos [11, 12].

$$\nabla_\mu J^\mu_L = \frac{N_{R-L}}{24\pi^2} R\tilde{R}$$

(1)

where $J^\mu_L$ is the total lepton current and $N_{R-L}$ is the difference between the number of right-handed and left-handed lepton species. As a matter of fact, it has been shown that neutrino masses, either Dirac or Majorana, do not affect the predictions of the gravitational leptogenesis [13].

The necessary ingredient in this case for the generation of a net lepton number is the existence of a primordial chiral gravitational wave background which contributes to the Pontryagin density $R\tilde{R}$. In order to generate such a chiral background, extended inflationary models involving axial couplings of the inflaton field have been considered. Thus for example, a gravitational Chern-Simons coupling of a pseudo-scalar inflaton field was originally proposed in [10], although some consistency issues were discussed in [14, 15]. Other possibilities include a Chern-Simons interaction between the pseudo-scalar inflaton and a $U(1)$ gauge field [16] and non-abelian gauge fields coupled to an axionic inflaton [17, 18]. Alternative ways of generating a parity-violating GW background have been considered in [19–21].

In this work we make the observation that the gravitational leptogenesis mechanism can be implemented without invoking new axial couplings in the inflaton sector. Indeed, the perturbed Robertson-Walker (RW) background emerging after inflation already breaks parity thus generating a non-vanishing Pontryagin density. Notice that although the probability distribution functions for the production of left and right handed gravity waves are the same in ordinary inflation models, our universe is a particular realization of the Gaussian process in which the actual amplitude of left and right handed gravitational wave excitations can be different.

We thus conclude that the minimal Standard Model with left-handed neutrinos together with an ordinary inflationary model driven by a scalar inflaton field already contains all the ingredients to generate a lepton asymmetry after inflation.

II. GRAVITATIONAL LEPTOGENESIS

Let us then consider a spatially flat RW spacetime with scalar and tensor perturbations in the longitudinal gauge. We will ignore vector perturbations as they are not typically produced during inflation. The line element in conformal time reads

$$ds^2 = a^2(\eta) \left[ (1 + 2\Phi) d\eta^2 - ((1 - 2\Psi)\delta_{ij} - h_{ij}) dx^i dx^j \right],$$

(2)

with $\Phi, \Psi$ the scalar perturbations and $h_{ij}$ the transverse traceless tensor perturbation.
The Pontryagin density that sources the leptonic current in (1) can be written in terms of the electric and magnetic parts of the Weyl tensor \([22–24]\), \(E\) and \(B\) respectively, as
\[
R \tilde{R} = \frac{1}{2} \varepsilon_{\mu \rho \sigma} R_{\mu \nu \alpha \beta} R^{\rho \sigma \alpha \beta} = 16 E_{\mu \nu} B^{\mu \nu}.
\]

The unperturbed part of the metric (2) does not contribute to either the electric or magnetic parts of the Weyl tensor. On the other hand, all scalar, vector and tensor components contribute to the electric part, whereas only vector and tensor ones add to the magnetic part. This means that the leading contribution to \(R \tilde{R}\) is second order in metric perturbations. Since the tensor-to-scalar ratio of the primordial power spectra \(r < 0.1\) \([25]\), we expect the scalar-tensor contribution to dominate over the tensor-tensor one, i.e.
\[
R \tilde{R} = - \frac{4}{a^4} \epsilon_{ijkl} (\Phi + \Psi)_{ij} h^i_{, k l} + \ldots
\]
where prime denotes derivative with respect to the conformal time. In the comoving frame we can write \(J^k_L = (a^{-1} n_L, 0)\) to leading order in perturbations, where \(n_L\) is the physical lepton number density. Inserting these expressions into (1) we obtain the leptonic number density after integrating in time
\[
n_L = \frac{1}{16 \pi^2 a^3} \epsilon_{ijkl} \int \mathrm{d} \eta \Phi_{,ij} h^i_{, k l},
\]
where we used that in the Standard Model \(N_{R-L} = -3\), and that in the absence of anisotropic stress \(\Phi = \Psi\).

Let us now expand the scalar and tensor perturbations in terms of creation and annihilation operators
\[
\Phi(\eta, \mathbf{x}) = \int \frac{\mathrm{d}^3 p}{(2\pi)^{3/2}} \left( \Phi(p, \eta) b_p e^{i \mathbf{p} \cdot \mathbf{x}} + \Phi^*(p, \eta) b_p^\dagger e^{-i \mathbf{p} \cdot \mathbf{x}} \right)
\]
and
\[
h_{ij}(\eta, \mathbf{x}) = \int \frac{\mathrm{d}^3 k}{(2\pi)^{3/2}} \sum_{\lambda = +, -} \left( h_\lambda(k, \eta) e^{i\lambda}_j(k) \alpha_{k, \lambda} e^{i k \cdot \mathbf{x}} + h_\lambda^*(k, \eta) e^{i\lambda^*}_j(k) \alpha_{k, \lambda}^\dagger e^{-i k \cdot \mathbf{x}} \right),
\]
where \(\lambda = \pm\) correspond to the \(\pm 2\) helicity modes whose polarization tensors can be written as \(e^{i\lambda}_j(k) = e^{i\lambda^*}_j(k) \varepsilon^\lambda(k)\), with \(\varepsilon^\lambda(k)\) the helicity \(\pm 1\) polarization vectors. Notice that in the absence of chiral couplings \(h_+ = h_- = h\).

It is straightforward to see that the expectation value of the scalar-tensor contribution to the lepton number in the Bunch-Davies vacuum is vanishing. Indeed, we can schematically write \(\langle n_L^2 \rangle \sim \langle \Phi h^* \rangle = \langle \Phi \rangle h^* = 0\). Therefore, the leading contribution to the expectation value would be the tensor-tensor one, which has been already explored in previous works \([6, 10, 16]\). However this contribution also vanishes for inflationary sectors without axial couplings. Notice however that the variance of the lepton number density is in general non-vanishing since we can write \(\langle n_L^2 \rangle \sim \langle \Phi^2 \rangle \neq 0\). Precisely the root mean square \(n_{L, \text{rms}}^2 = \langle n_L^2 \rangle^{1/2}\) provides an estimate of the produced lepton density in a typical realization of the random process. Thus, it is straightforward to obtain
\[
\langle n_L^2 \rangle = \left( \frac{1}{16 \pi^2 a^3} \right)^2 \sum_{\lambda \in \{+,-\}} \int \frac{\mathrm{d}^3 k \, \mathrm{d}^3 p \, \mathrm{d} \eta \, \mathrm{d} \eta'}{2 \pi^6} \times k^2 |p \cdot e^{i\lambda}_j(k)|^4 \langle h^*_\lambda(k, \eta) h^{*\lambda}(k, \eta') \Phi(p, \eta) \Phi^*(p, \eta') \rangle
\]
In the absence of chiral couplings the whole integral in (8) is independent of \(\lambda\) thanks to spherical symmetry. After some simplification, the variance can be written in a compact manner as
\[
\langle n_L^2 \rangle = \frac{1}{384 \pi^2 a^6} \int \mathrm{d} k \, \mathrm{d} p \, k^4 p^6 \int \mathrm{d} \eta \Phi(p, \eta) h^*_I(k, \eta) \bigg| \bigg| h^*_I (k, \eta) \bigg| \bigg|^2
\]

### III. LEPTOGENESIS DURING REHEATING

For the sake of concreteness, we will assume that the net lepton number density at the end of inflation is negligible, so that we will consider the leptogenesis produced throughout the stage of reheating by the inflationary primordial metric perturbations. For simplicity, we consider that the energy content during reheating is described by means of an effective fluid with barotropic equation of state \(p = w \rho\), with \(w\) constant. We will also parametrize the primordial power spectra in the usual way:
\[
\mathcal{P}_S(k) = A_S \left( \frac{k}{k_x} \right)^{n_s - 1}, \quad \mathcal{P}_T(k) = A_T \left( \frac{k}{k_z} \right)^{n_T},
\]
where \(A_S\) and \(A_T\) are the scalar and tensor amplitudes at the pivot scale \(k_x\), respectively. We will use the value obtained by the Planck collaboration \([2]\) for the scalar spectral index \(n_s = 0.965\) and assume a scale-invariant tensor power spectrum \(n_T = 0\). The primordial power spectra generated during inflation have a natural ultraviolet cutoff at the scale \(k_I = a_t H_I\) corresponding to the size of the comoving Hubble horizon at the end of inflation, as modes with \(k > k_I\) have never left the horizon and could not become classical metric perturbations. Imposing this upper limit in the momentum integrals in (9), we obtain for the total lepton number variance generated during reheating
\[
\langle n_L^2 \rangle_{\text{RH}} = \frac{1}{960 \pi^4} \left( \frac{k_I}{a_{\text{RH}}} \right)^6 \mathcal{P}_S(k_I) \mathcal{P}_T(k_I) I(\eta_{\text{RH}})
\]
with
\[
I(\eta) = \int_0^1 \mathrm{d} x \int_0^1 \mathrm{d} y \, x^{n_T + 1} y^{n_s + 2} \times \left| \int_{\eta_L}^{\eta} \mathrm{d} \eta' \tilde{\Phi}(k_I x, \eta) \tilde{h}^*(k_I x, \eta) \right|^2
\]
which correspond to modes with $k$ integrals are dominated by the upper integration limits $x$ similarly for the tensor mode. Notice that the $x$ and $y$ integrals are dominated by the upper integration limits which correspond to modes with $k \approx p \approx k_f$.

During reheating, scalar and tensor modes behave as

$$\Phi(p, \eta) = \eta^{-r}[C_1 J_r(\sqrt{w} p \eta) + C_2 Y_r(\sqrt{w} p \eta)]$$

and

$$h(k, \eta) = \eta^r[D_1 J_s(k \eta) + D_2 Y_s(k \eta)]$$

with $J_{r,s}$ and $Y_{r,s}$ Bessel functions of order $r = (5 + 3w)/(2 + 6w)$ and $s = (3w - 3)/(2 + 6w)$ respectively and $C_{1,2}$, $D_{1,2}$ constants. Both scalar and tensor modes remain constant outside the horizon ($k \eta \ll 1$), while well inside the horizon they oscillate with decreasing amplitude, except for $w = 0$, where $\Phi = \text{const.}$ even inside the horizon. In Fig. 1 we show the behaviour of the perturbations together with the product appearing in the integrand of (12) for different values of $w$ for the dominant modes. We can see that contributions to the time integral in (12) occur only when tensor modes enter the Hubble horizon, since in the super-Hubble regime $h' \approx 0$. We also plot function $I(\eta)$ in Fig. 2 which provides the time dependence of the lepton production. The different behaviour for each equation of state, and hence the different interference between scalar and tensor modes, translates into a different time evolution of $I(\eta)$. We can also see that leptogenesis takes place in a few Hubble times.

We can now compute the lepton number to entropy ratio by dividing by the entropy density $s = 2\pi^2 g_s T_{RH}^4/45$, with $g_s$, the effective number of relativistic species. This ratio should be so that it reproduces the baryon asymmetry after the partial conversion of leptonic asymmetry via sphalerons [4], which is

$$\frac{n_L^{\text{rms}}}{s} = \frac{79}{28} \left| \frac{n_B}{s} \right| = 2.45 \times 10^{-10}. \quad (15)$$

where $a_{RH}$ denotes the scale factor at the end of reheating, $\Phi(k, \eta) = \Phi(k, \eta)/\Phi(k, \eta_f)$ is the scalar perturbation normalised to its value at the end of inflation, and similarly for the tensor mode. Notice that the $x$ and $y$ integrals are dominated by the upper integration limits which correspond to modes with $k \approx p \approx k_f$.

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Far from the instantaneous reheating limit we find that in general $I(\eta_{RH}) = O(10^{-1})$. Thus, from (11), we see that, apart from the power spectra which are (almost) scale invariant, the relevant quantity in determining the size of $n_L^{\text{rms}}/s$ is the factor $(k_f/a_{RH})^3$. As a matter of fact, we can extract the dependence on $H_I$ and $T_{RH}$ from such a factor and obtain

$$\frac{n_L^{\text{rms}}}{s} \left| \frac{H_I}{T_{RH}} \right| \approx \sqrt{A S A T} \left( \frac{H_I}{M_P} \right)^{\frac{1+3w}{1+w}} \left( \frac{T_{RH}}{M_P} \right)^{\frac{1-3w}{1+w}} \quad (16)$$

for $n_s \approx 1$ and where $M_P = 1/\sqrt{G}$ is the Planck mass, meaning that leptogenesis is enhanced for high inflation scales and, in the case of stiff reheating scenarios with $w > 1/3$, for low reheating temperatures.

In the left panel of Fig. 3, we plot $n_L^{\text{rms}}/s$ in the $(T_{RH}, H_I)$ parameter space for a stiff equation of state $w = 1$ during reheating. We see that the asymmetry in
of state during reheating

and tensor spectra, it is possible to obtain analytical ex-

parameters used in that work, we get

we find that for the instantaneous reheating case and the

results with those obtained in [10] with the axial coupling,

sufficient at producing lepton number. Comparing these re-

w = 1/3, we find that, as expected, the production is not

sensitive to the reheating temperature. Reheating sce-

narios with equations of state close to w = 0 are not effi-

cient at producing lepton number. Comparing these re-

results with those obtained in [10] with the axial coupling,

we find that for the instantaneous reheating case and the

parameters used in that work, we get \( n_L^{\text{rms}}/s \sim 10^{-20} \) which is slightly above their results.

IV. SIZE OF MATTER-ANTIMATTER REGIONS

The variance of the lepton number density obtained

in (11) only provides the typical amplitude of local fluc-

tuations. In order to determine the size of the matter-

antimatter regions it is necessary to calculate the corre-

lation function \( \xi(r) = \langle n_L(x+r)n_L(x) \rangle \). For an equation

of state during reheating \( w = 0 \) and scale invariant scalar

and tensor spectra, it is possible to obtain analytical ex-

pressions. Thus, changing variables \( p = q - k \), we can write

\[
\xi(r) = \frac{A_S A_T}{4096 \pi^4 a^b} \int_{k_0}^{k_I} dk \int d^3q \, k q^4 e^{i q r} \times \left( \frac{1 - x^2}{q^2 + k^2 - 2 k q x} \right)^{1/2},
\]

where \( k_0 \) and \( k_I \) denote the infrared and ultraviolet cut-

offs of the production. The following limits can be obtained

\[
\xi(r) = \begin{cases} 
\frac{41 A_S A_T}{483840 \pi^4} \left( \frac{k_I}{a} \right)^6, & k_I r \ll 1 \\
\frac{11 A_S A_T}{26880 \pi^4} \left( \frac{k_I}{a} \right)^6 \cos( k_I r) \frac{1}{k_I r^2}, & k_I r \gg 1 
\end{cases}
\]

The comoving coherence length associated to the lep-

ton number variance is therefore \( l_{\text{coh}} \approx 1/k_I \), which is

much smaller than the observable universe. We thus ex-

pect that these small matter-antimatter regions will an-

nihilate each other on subhorizon scales, providing an

average asymmetry over a comoving scale \( r_0 \) given by

the weighted variance

\[
\langle n_L^2 \rangle_{r_0} = \left( \frac{1}{V_W} \left( \int d^3r n_L(x+r) W(r) \right) \right)^2,
\]

where \( W(r) \) is a window function of characteristic size \( r_0 \) and \( V_W = \int d^3r W(r) \). Again for an equation of state
during reheating $w = 0$ and scale-invariant spectra we can write

$$
\langle n_L^2 \rangle_{r_0} = \frac{A_s A_T \pi}{256 \alpha^6 r_0} \int_0^{k_f} dk \int_0^{k_i} dq q^6 |W(k + p)|^2 V_W^2 \\
\times \int_{-1}^{1} dx \left( \frac{1 - x^2}{q^2 + k^2 - 2kqxe^{3/2}} \right). \tag{20}
$$

If we choose a gaussian window function $W(r) = e^{-r^2/2r_0^2}$, we can obtain

$$
\langle n_L^2 \rangle_{r_0} = \frac{A_s A_T}{2048 \pi^4 a^6 r_0} \left[ \frac{2e^{-k^2 r_0^2}}{k_0^2 r_0^2} + \sqrt{\pi} \left( \frac{1}{k_f^2 r_0^2} - \frac{2}{k_1 r_0} \right) \right] \\
\times \text{erf}(k_f r_0) - \left( \frac{1}{k_0^2 r_0^2} - \frac{2}{k_0 r_0} \right) \text{erf}(k_0 r_0). \tag{21}
$$

Compared to the local variance in (8), this quantity is suppressed as

$$
\langle n_L^2 \rangle_{r_0} \simeq \frac{1}{(k_f r_0)^6} \frac{1}{k_0 r_0} \langle n_L^2 \rangle, \tag{22}
$$

exhibiting a clear blue-tilted behaviour. Thus, on regions of order $r_0 = H_0^{-1}$, the expected lepton asymmetry will be several orders of magnitude smaller than observations.

V. PHENOMENOLOGY IN THE LATE UNIVERSE

Let us examine the implications of this gravitational leptogenesis mechanism for standard cosmology. Firstly, the lepton number asymmetry generated during reheating can be converted into baryon asymmetry only if $T_{RH}$ is above the electroweak scale. Around and above this temperature, QCD confinement has not occurred yet, so the baryon number is in the form of quarks, which are relativistic. Quarks interact in this pre-confinement plasma with a mean free path which can be estimated as $\Gamma_q^{-1} \simeq T^{-1}$, causing diffusion of the baryon number. As a result, perturbations in the baryon asymmetry are suppressed on scales below the corresponding (comoving) Silk length which can be estimated at the moment of confinement as follows [26]

$$
r_S^2(T_{QCD}) = \int_0^{a_{QCD}} \frac{\Gamma_q^{-1}(a)}{a^3 H(a)} \simeq (10^{-16} \text{ Mpc})^2, \tag{23}
$$

with $a_{QCD}$ the scale factor at the confinement temperature $T_{QCD} \simeq 300$ MeV.

After confinement, quarks cannot longer exist as free particles and form bound states, namely protons and neutrons. These particles are now non-relativistic and still interact with photons, which makes the baryon diffusion scale drop significantly, so the comoving size of the matter-antimatter patches freezes after confinement. We can calculate the weighted variance of the baryon asymmetry at the Silk scale $r_0 = r_S$ from (19). Note that, after inserting (11) into (22), the dependence on the inflation scale cancels out (except for the power spectra, which are nearly scale-invariant), so the Silk scale becomes the only relevant one so that we obtain

$$
\langle n_B^2 \rangle_{rs}^{1/2} \simeq 10^{-36} \left( \frac{H_0}{k_0} \right)^{1/2}, \tag{24}
$$

for an infrared cutoff $k_0$.

Thus we have the following behaviour for the $rms$ baryon asymmetry fluctuations at a given $r_0$ scale

$$
\frac{n_B^{rms}}{s} \simeq \begin{cases} 
10^{-36} \left( \frac{H_0}{k_0} \right)^{1/2}, & r_0 < r_S \\
10^{-36} \left( \frac{H_0}{k_0} \right)^{7/2}, & r_0 > r_S
\end{cases}, \tag{25}
$$

i.e. for regions smaller than the Silk scale, diffusion suppresses baryon fluctuations and the abundance must be obtained through the weighted variance over the Silk length, which does not depend on the inflationary parameters. For larger patches $r_0 > r_S$, the averaged abundance is damped as shown in Eq. (22).

In conclusion, we see that gravitational leptogenesis effects will induce tiny fluctuations in the baryon asymmetry parameter over regions with a typical size of the Silk length at confinement. Regarding the lepton number asymmetry that is also generated, it is homogenised in a similar manner through free-streaming of electrons and especially neutrinos, which are relativistic until very late stages of cosmic evolution.

VI. DISCUSSION

In the previous analysis we have considered simple power laws for the scalar and tensor primordial spectra in the whole range of scales with the amplitudes and spectral indices measured from CMB observations. However, as shown in (11), the produced lepton density depends on the values of the power spectra at the $k_f$ scale, which can be separated from the scales measured in the CMB by many orders of magnitude. This means that a possible running of the spectral indices could affect the predictions of the model. An interesting possibility would be the presence of features in the scalar power spectrum at small scales. In particular, it has been shown that the presence of (broad) peaks could play an important role in the generation of primordial black holes after inflation [27]. Such peaks could in fact enhance in several orders of magnitude the produced lepton asymmetry.

As shown in Fig. 3, the lepton to entropy ratio strongly depends on the equation of state during the reheating phase and stiff equations enhance the generated asymmetry. This possibility has been recently discussed in [6] where kination dominated reheating scenarios [28]
have been considered. An interesting consequence of early phases with stiff equation of state is the generation of a blue tilt in the transfer function of tensor modes [29] which could render the primordial gravitational wave background observable for the sensitivity and frequency range of future detectors such as LISA, Einstein Telescope or Cosmic Explorer. Even when considering the rms fluctuation at Silk scale, which eliminates the explicit dependency on the inflation scale, both a kination phase and the presence of features in the power spectra introduce a dependency on the particular inflationary scenario.

Even though primordial metric perturbations do not seem to be able to generate the observed homogeneous asymmetry on Hubble scales, the produced baryon asymmetry could in principle act as a source of baryonic isocurvature perturbations. However, according to the obtained results (25), these perturbations would be very small for observable scales. Finally, let us mention that beyond the linear regime, gravitational lepton generation in chiral astrophysical systems [22] could also provide potential experimental ways to test the leptogenesis mechanism discussed in this work.

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