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Implications of Non Standard Scenarios in Cosmology and the Very Early Universe

PhD dissertation by

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ABBREVIATIONS

This a list of the abbreviations used throughout this manuscript:

RD Radiation domination

MD Matter domination

FLRW Friedman-Lemmaitre-Robertwon-Walker

BBN Big Bang Nucleosynthesis

SM Standard Model

EWPT Electroweak phase transition

LHC Large Hadron Collider

DM Dark Matter

CDM Cold Dark Matter

CMB Cosmic Microwave Background

GUT Grand Unification Theories

RH Right handed

VEV Vacuum expectation value

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Part I

GENERAL INTRODUCTION

The theoretical picture of the Universe nowadays is widely accepted to rest on the Standard Model of the fundamental interactions and the Big Bang Cosmological model. The former one is responsible for the description of Nature at the particle size scale, whereas the latter model gives rise to a picture for the dynamics and evolution of the Universe at the macroscopic level. It is such the success that one usually refers to both theories as the standard point from which any extending model should be formulated.

A look into the macroscopic level

Cosmology is responsible for the study of the origin, evolution and fate of the Universe. At this macroscopic level, the dynamics of the Universe can be addressed by invoking general relativity. As a consequence, any cosmological model is defined by the components that fill the Universe and the geometry of space-time.

As usual, the construction of a particular physical model is inspired by evidences. Regarding the Universe as a whole, the expansion of the Universe and its isotropy and homogeneity on large scales are the observational cornerstones. Both evidences led to the formulation of the Big Bang model, which is considered as the Standard model in Cosmology. Such a model is based on the Friedmann-Robertson-Walker metric, which, in combination with the particle content as a perfect fluid, provides a suitable description of the Universe. According to this, the Universe was born from a singularity 13 billion years ago and it began to rapidly expand. As this expansion came along, the Universe cooled down sufficiently to allow the formation of the first atomic nuclei and light elements, an important moment in the History of the Universe known as the Big Bang Nucleosynthesis. It is important to point out that the concordance with BBN is so high that, on the theoretical side, it is one of the most stringent probes of the Standard Scenario triggered by the Big Bang model.

On the other hand, the Big Bang Cosmological Model also provides with an ideal scenario for the generation of the observed astronomical objects through a process of gravitational baryonic attraction once the expansion of the Universe sufficiently slows down at late times. The only ingredient assumed for this picture to be viable is the existence of a weakly interactive non relativistic matter called “Dark Matter”, which dominated the Universe some time after BBN. The existence of this type of matter was first confirmed on the galaxy rotation curves and measured to be nowadays the 26% of the Universe fluid.

Finally, the standard picture is completed with the inclusion of a component called “dark energy”, which is responsible for the today accelerated expansion of the Universe and fills the 70% of the current total energy of the Universe

In sum, the FLRW metric along with the fluid decomposition into radiation, non relativistic matter (Dark and baryonic matter) and Dark Energy set up our standard understanding of the Universe as a whole.

Nevertheless, such a standard scenario is far from being complete. One of most prominent issues is related to the homogeneity and isotropy paradigm itself which, given what we know about the standard thermal history of the Universe, could have never been reached within the Big Bang scenario. The most accepted explanation to solve this is to include a brief period of exponential expansion of the Universe after the Big Bang known as Inflation. It is so widely accepted that most of the scientific community includes it as a standard ingredient yet its nature and the definitive Inflationary model is still unknown.

Moreover, Inflation is commonly assumed to be connected to the BBN by a long period of radiation domination. However, there is no evidence for such an assumption and one can explore deviations to this standard scenario that can have important consequences. Therefore, it can be seen then that the whole puzzle is not completed yet and extensions to the Big Bang picture need to be incorporated or at least deserve being investigated.

A look into the microscopical level

The Standard Model of particle physics elegantly provides a theoretical framework for the underlying phenomenology below energy scales of 100 GeV by simply invoking to quantum field theory and symmetry arguments. The symmetry, usually known as a *gauge* symmetry, that governs the allowed operators which may appear in the Lagrangian for the Standard Model is $SU(3)_c \otimes SU(2)_L \otimes U(1)_Y$, where “c” is the colour charge, “L” the left-handed isospin and “Y” the Hypercharge.

Such operators involve quantum fields that are arranged according to their quantum numbers, where each quantum field involves the physical particle content organised in suitable representations. In this way, Quarks are represented by the multiplets $\mathcal{Q}$, u and d, with these last fields being singlets under $SU(2)_L$. Leptons, on the other hand, are represented by $\mathcal{L}$ and e_R, both singlets under $SU(3)_c$ and e_R under $SU(2)_L$ too. The interactions among these fields are given by the gauge bosons, as the interaction carriers of the fundamental forces, that naturally arise in the theory as a consequence of the gauge invariance.

The success of the Standard Model of particle physics is unquestionable. It did not only account for the observed phenomena that take place in nature, but also made predictions that have been confirmed experimentally. More recently, the last prediction, as being the leftover part of the physics description at the $\mathcal{O}(100GeV)$ scale, the Higgs boson, responsible for the origin of the particle masses, has been discovered at the LHC.

Even though its great success, many observational evidences remain a mystery when one attempts to resort to the Standard Model to explain them. Among all these unsolved problems, I would like to highlight three of them: the origin of the baryon asymmetry, the nature of dark matter and the neutrino masses.

The first open question has to do with the fact that we observe almost no track of antimatter, i.e, the observable Universe is mainly composed entirely of ordinary matter. In principle, the Standard Model has all the needed ingredients to generate an asymmetry between matter and antimatter, but it fails to do so due to the value of the parameters in play.

On the other hand, Dark matter is widely accepted to have the following features: being weakly interactive with the rest of particles and cold *i.e*, non relativistic, when it decoupled from the thermal bath. As a consequence of this, it is easy to see that there is no plausible candidate to be the Dark Matter within the Standard Model given its particle content.

The discovery of the Higgs boson, as I previously mentioned, appears to be the explanation for the masses of the particles. However, some time ago, it was found that the neutrinos also have mass. This feature, even with the inclusion of the Higgs boson, can not be addressed in the SM unless one resorts to a particle extension and/or the introduction of a new physics (Neutrino Dirac mass terms do not need new physics).

In sum, it is then clear that the picture at the microscopical level too is still unclear and incomplete, so extensions to the scenario proposed by the SM need to be taken into account.

A non standard approach

As it has been made clear in the previous sections, new extensions to the standard scenarios are welcome. In this Thesis then, the features of some of these scenarios have been investigated with the aim of providing some insight in the open and unsolved questions and look for the observational traces that these new scenarios can leave. In order to do so, the approach has been twofold: Right Handed Neutrino composites in the very early stages of the Universe and Matter Domination prior to BBN.

The existence of scalar particles in Nature has always disturbed the physics community due to the issues that arise, specially the natural introduction of hierarchies in the theory. As a consequence, many models have been proposed where scalar particles are no longer elementary but a composite of fermions, which lack of hierarchy issues. Most of this kind of works have been focused on the nature of the Higgs particle such as in technicolour and extended versions [1–6]. There are other models where even gauge bosons can naturally consist of composites of fermions [7]. In all these theories, the hierarchy issues are solved by the introduction of new physics where the electroweak scale arises naturally. Moreover, cancellations are no longer needed to keep the electroweak scale under control.

We have thus wondered whether the Inflaton as being a scalar particle can be explained in the same way. Measurements set the inflationary scale to be around 10^{16}GeV, whose origin might be in principle quite controversial. We then explored the origin of such a scale to come from a composite of majorana RH neutrinos, since as expected to be very heavy in order to account for the size of the observe neutrino masses by the seesaw mechanism, one could reason that both scales are related. One then introduces an effective 4-Fermi operator, made up of RH neutrinos, with a cutoff scale Λ below which neutrinos can condense into a scalar particle if the so-called gap equation is satisfied. Such a scalar particle plays the role of the Inflaton, breaking down the Lepton number symmetry spontaneously when it develops a vacuum expectation value. Nevertheless, such a mechanism by itself does not provide an inflationary potential, the composite Inflaton would be massless and hence inflation could not be possible. However, due to the lepton number violation from the inclusion of Majorana RH neutrinos, the symmetry is broken explicitly and the composite arises now as a Pseudo Nambu-Golstone boson, with a potential of the form $V(\phi) = M^4(1 + \cos(\phi/\mu))$.

Furthermore, it is easy to show that the baryon asymmetry can be naturally generated within this model with the mentioned minimal inclusion of right handed neutrinos in the theory. The Inflaton as a composite, can derivatively couple to the lepton current if CPT is violated, inducing then a different chemical potential for baryons and anti-baryons. These quantities will then trigger a net lepton number, which transformed into the observed baryon asymmetry through the sphalerons at EWPT.

On the cosmological side, we have also turned our attention to the behaviour of the Universe between Inflation and BBN. As it was demonstrated in [8], a period of Matter domination in between can play a somewhat similar role to Thermal Inflation [9] without the inclusion of a scalar particle. We continue this line by studying the effects that this early period of matter can leave on the electroweak phase transition and the formation of structures.

One of the reasons that makes the SM fail to give rise to the baryon asymmetry is that the electroweak phase transition is not first order enough. Whether the electroweak phase transition is strong enough is determined by the so-called sphaleron bound, which relates the Higgs vacuum expectation value $\langle\phi\rangle$ and the critical temperature T_c. It turns out that if the baryon asymmetry is to survive after the electroweak phase transition, the sphaleron bound states that $\langle\phi\rangle > T_c$, a bound which can be translated into a constraint on the parameters of the theory. It is easy to demonstrate that the SM does not have enough number of bosons with the right spectrum to satisfy this bound. The common alternative to overcome this issue is the increase of the number of bosons in the theory, that raise then the strength of the EWPT. However, our approach is different and we show that under a different thermal history of the Universe with the inclusion of an early matter domination epoch, the sphaleron bound can be relaxed and the window for the baryon asymmetry generation in the standard scenarios reopen.

Finally, an intermediate matter domination stage can also have consequences on the evolution of the primordial perturbations. As it is known, perturbations inside the horizon grow during MD much faster than in radiation domination. In the standard picture, such a growth, which leads to the objects that we observe, takes place when Dark Matter came to dominate the Universe long time after BBN. But this raises the question: What if this growth takes place before BBN due to an early period of matter domination?. We consider this scenario, that can be fully characterised by the features of the particle(s) that dominate(s) the Universe during that early stage. We then make use and evolve the equations of motions for perturbations during this epoch and show that a matter dominated Universe before BBN can leave traces in the form of structures and substructures that in principle can be detected.

The Thesis structure

This thesis aims to explore extensions to the standard scenario which have been considered in the literature so far. It consists of a collection of some of my published papers during my PhD. They have been selected in order to follow a coherent topic exposition in this document.

The parts of the Thesis following this chapter are the following:

- A part to introduce and review the minimal concepts in the Literature which might be useful to understand the papers. It has been divided into 4 chapters. In the first chapter the Standard Cosmological Model and the standard history of the Universe are briefly introduced outlining its main features. Inflationary physics is then explained with great detail in a new chapter. It then follows a chapter dedicated to the generation of the baryon asymmetry in the SM. This part concludes with a thorough review of the standard structure formation picture.

- In the following part, my research contribution is presented by the inclusion of the papers where I have taken part. In the first chapter, the Inflationary and baryogenesis scenario from a condensate of Heavy Right Handed neutrinos is explained. The following chapter studies the electroweak phase transition in the SM and extensions when an early period of matter domination in the Universe is introduced. The final chapter of this part applies the same features of the former scenario but now to the growth and amplification of primordial perturbations.

- I conclude with a part dedicated to summarise the main results and the conclusions that can be extracted from the papers. Likewise, the literature used in the review part and publications is listed.

INTRODUCCIÓN GENERAL

Está ampliamente aceptado hoy en día que el marco teórico del Universo está formado por el Modelo Estándar de las interacciones fundamentales y el modelo cosmológico del Big Bang. El primero es responsable de la descripción de la naturaleza a la escala del tamaño de las partículas, mientras que el segundo da lugar a la dinámica y evolución del Universo a nivel macroscópico. Es tal el éxito de ambos modelos que generalmente sirven como punto de partida para cualquier extensión teórica.

La Cosmología se encarga del estudio del origen, evolución y destino del Universo. A este nivel macroscópico, la dinámica del Universo puede ser explicada mediante Relatividad General, lo que hace que cualquier modelo cosmológico quede definido por las componentes del Universo y su geometría del espacio-tiempo.

Como es costumbre, la construcción de cualquier modelo físico se encuentra promovido por ciertas evidencias observacionales. En el caso del Universo en su conjunto, la expansión de éste y su isotropía y homogeneidad a grandes escalas suponen sus pilares básicos. Ambas evidencias dan lugar al Modelo del Big Bang, que se puede considerar como el modelo estándar en Cosmología. Dicho modelo está basado en la métrica de Friedmann-Robertson-Walker que, combinándola con la consideración del contenido de partículas en fluidos perfectos, proporciona una descripción ajustada del Universo. De acuerdo a esto, el Universo nació de una singularidad hace 13.000 millones de años y empezó entonces a expandirse rápidamente. A medida que se expandía, el Universo se enfrió lo suficiente para permitir la formación de los primeros núcleos atómicos y elementos ligeros, hechos que se produjeron en el instante de la historia del Universo conocido como Big Bang Nucleosíntesis. Es importante aclarar que es tal el grado de concordancia del modelo del Big Bang con dicho momento que supone una de las pruebas más robustas del escenario estándar.

Por otro lado, el modelo cosmológico del Big Bang también proporciona, mediante un proceso de atracción gravitatoria de materia bariónica, un escenario ideal para la generación de los objetos astronómicos que observamos hoy en día. El único ingrediente que se asume para hacer este marco viable es la existencia de una componente de materia no relativista conocida como “Materia Oscura”, que apenas interactúa con el resto de partículas y que dominó el Universo algún tiempo después de BBN. La existencia de este tipo de materia fue confirmada por primera vez gracias a la observación de las curvas de rotación de las galaxias y se ha medido recientemente que supone el 26 % del fluido total del Universo. Finalmente, el marco estándar se completa con la inclusión de un componente llamado “Energía Oscura” que es el responsable de la expansión actual acelerada del Universo y que llena el 70% del mismo.

Para concluir, la métrica de FLRW junto con la descomposición del fluido del Universo en radiación, materia no relativista (Materia Oscura y bariónica) y Energía Oscura establecen nuestro entendimiento estándar del Universo en su conjunto.

No obstante, dicho escenario que consideramos estándar está aún lejos de estar completo. Uno de los problemas más notables tiene que ver con el mismo paradigma de que el Universo es homogéneo e isótropo ya que, dado lo que conocemos sobre la historia térmica del Universo, tales características no se podrían haber alcanzado en el escenario del Big Bang. La explicación más aceptada para resolver esto consiste en la incorporación de un breve periodo de expansión exponencial del Universo justo después del Big Bang conocido como Inflación. Dicho periodo es tan ampliamente aceptado dentro de la comunidad científica que se puede considerar como una componente más dentro del marco teórico estándar aunque su naturaleza y modelo definitivo aún no se hayan encontrado.

Además, se suele considerar que inflación está conectado con BBN por un largo periodo de dominación de radiación. Sin embargo, esto es simplemente una asunción sin ninguna evidencia que la confirme así que uno puede explorar alternativas a este escenario estándar y que pueden tener importantes consecuencias. Por tanto, se puede concluir que el rompecabezas general aún no está completo del todo y extensiones al marco generado por el Big Bang tienen que ser incorporadas o al menos merecen que se investiguen.

Una mirada al mundo microscópico.

El modelo Estándar de física de partículas proporciona de manera elegante un marco teórico de la fenomenología por debajo de los 100

GeV's simplemente usando la Teoría cuántica de campos e invocando a argumentos de simetría. La simetría en la que se basa el Modelo Estándar (conocida comúnmente como de gauge) y que determina los operadores que aparecen en su Lagrangiano es $SU(3)_c \times SU(2)_L \times U(1)_Y$, done "c" es la carga de color, "L" el isospin de levógiro e "Y" la hipercarga. Estos operadores mencionados involucran campos cuánticos que están dispuestos de acuerdo a sus números cuánticos, donde cada campo cuántico contiene las partículas físicas organizadas mediante representaciones adecuadas. De esta manera, los quarks vienen representados por los multipletes $\mathcal{Q}$, u y d, siendo estos dos últimos singletes bajo $SU(2)_L$. Los leptones, por otro lado, vienen representado por $\mathcal{L}$ y e_R, ambos singletes bajo $SU(3)$ y e_R también bajo $SU(2)_L$. Las interacciones entre estos campos vienen dadas por los bosones de gauge, que actúan como portadores de las fuerzas y que aparecen de manera natural debido a la invariancia gauge.

El éxito del Modelo Estándar de partículas está fuera de toda duda. No sólo fue capaz en su momento de explicar los fenómenos que se observaban sino que dio lugar a nuevas predicciones que han podido ser confirmadas tiempo después. La última predicción del Modelo Estándar es el bosón de Higgs, responsable del origen de la masa de las partículas y que ha sido recientemente descubierto en el LHC.

Sin embargo, a pesar de este gran éxito, aún existen evidencias observaciones que el Modelo Estándar no ha sido capaz de explicar. Entre ellas, me gustaría resaltar las siguientes: el origen de la asimetría bariónica,la naturaleza de la materia oscura y la masa de los neutrinos.

La primera cuestión tiene que ver con el hecho de que apenas se encuentra antimateria en el Universo. En principio, el Modelo Estándar posee todos los ingredientes necesarios para explicar esta asimetría entre materia y antimateria, pero es incapaz de generarla debido a los valores de sus parámetros fundamentales.

Por otro lado, es ampliamente aceptado que la materia oscura tiene las siguientes características: interacciona débilmente con el resto de partículas y es fría, es decir, que era no relativista cuando se desacopló del baño térmico. Como consecuencia de esto, se puede ver fácilmente dado el contenido de partículas del Modelo Estándar que no existe un candidato plausible para ser materia oscura.

Como comenté anteriormente, el descubrimiento del bosón de Higgs parece ser la explicación de la masa de las partículas. Sin embargo, se vio hace tiempo que los neutrinos también tienen masa y esto no puede ser explicado por el Modelo Estándar (incluso con la simple existencia del HIggs) a no ser que uno extienda el contenido de partículas o recurra a nueva física (masas de tipo Dirac no necesitan nueva física por ejemplo).

En resumen, está también claro que la descripción del nivel microscópico todavía está incompleta, por lo que es necesario considerar escenarios que completen al Modelo Estándar.

Una estrategia no estándar

Como se ha comentado anteriormente, escenarios no convencionales que completen nuestro conocimiento estándar son bien recibidos. En esta Tesis, las consecuencias de algunos de estos escenarios han sido investigadas con el fin de proporcionar algún entendimiento en los problemas aún sin resolver y buscar las trazas observacionales que dichos escenarios no convencionales puedan dejar. De esta forma, dos escenarios han sido propuestos: existencia de compuestos de neutrinos dextrógiros en la época temprana del Universo y la dominación de materia no relativista antes de BBN.

La existencia de partículas escalares siempre ha intrigado a la comunidad científica debido a los problemas que surgen de ellas, como por ejemplo la introducción de jerarquías dentro de la teoría. Como consecuencia de esto, se han propuesto muchos modelos en los que las partículas escalares no son fundamentales sino un compuesto de fermiones, que no presentan problemas de jerarquías. La mayoría de estos trabajos se han centrado en el bosón de Higgs, como por ejemplo en technicolor y sus versiones extendidas [1–6]. También se han considerado modelos en los que incluso los bosones de gauge podían ser compuestos de fermiones [7]. En todas estas teorías, los problemas con las jerarquías se resuelven mediante la implantación de nueva física desde donde la escala electrodébil aparece de manera natural. Además, en este tipo de escenarios la magnitud de escala electrodébil está bajo control ante las correcciones radiativas.

En nuestro caso, nos hemos planteado si el Inflatón también puede ser explicado de esta manera. La escala inflacionaria, cuyo origen puede ser bastante controvertido, se ha medido que está alrededor de los 10^{16}GeV. Entonces, nosotros consideramos que el origen de esta escala pueda venir de la existencia de un compuesto de Neutrinos dextrógiros de tipo Majorana ya que, debido a su naturaleza pesada para explicar las masas de los neutrinos por el mecanismo de seesaw, uno podría pensar que las escalas están relacionadas entre ellas. Así, se introduce un operador efectivo que está formado por neutrinos dextrógiros y que es del tipo 4-Fermi, con la introducción de una escala límite Λ por debajo de la cual estos neutrinos condensan en una partícula escalar si se cumple la ecuación de gap. Esta partícula escalar cumple el papel de Inflatón, rompiendo así espontáneamente el número leptónico cuando desarrolla un valor esperado en el vacío. No obstante, este mecanismo no produce un potencial inflacionario y el inflatón entonces no tendría masa. Sim embargo, debido a que los neutrinos dextrógiros de tipo Majorana violan el número leptónico, la simetría también se rompe explícitamente y el compuesto ahora se comporta como un pseudo bosón de Nambu-Goldstone, con un potencial de la forma $V(\phi) = M^4(1 + \cos(\phi/\mu))$.

Además, se puede mostrar fácilmente que la asimetría bariónica puede surgir de manera natural en este modelo. El inflatón, como un un compuesto, puede acoplarse de forma derivativa a la corriente leptónica si CPT se viola, lo que induce un potencial químico diferente para bariones y antibariones. Estas cantidades producen entonces un numero leptónico neto que se puede transformar posteriormente a través de los esfalerones en la asimetría bariónica que observamos.

Del lado cosmológico, nos hemos fijado también en el comportamiento del Universo entre Inflación y BBN. Como se demostró en [8], un periodo de dominación de materia no relativista entre ambos puede jugar un papel similar al escenario descrito por la inflación térmica [9] sin la necesidad de introducir una partícula escalar. Nosotros entonces continuamos con esta línea de trabajo estudiando el efecto que un periodo de materia no relativista puede dejar en la transición de fase electrodébil y la formación de estructuras.

Una de las razones por las que el Modelo Estándar es incapaz de producir la asimetría bariónica tiene que ver con que la transición de fase electrodébil no es suficientemente de primer orden. Que dicha transición sea suficientemente fuerte depende del conocido límite de esfalerón, que relaciona el valor esperado en el vacío del Higgs $\langle\phi\rangle$ y la temperatura crítica T_c. Así, resulta que para que la asimetría bariónica sobreviva después de la transición de fase electrodébil, el límite de esfalerón dictamina que $\langle\phi\rangle > T_c$, un límite que puede ser transformado en una restricción en el valor de los parámetros de la teoría. Se puede demostrar fácilmente que el Modelo Estándar no tiene el suficiente número de bosones con el espectro correcto para satisfacer este límite. Entonces, la alternativa que se suele adoptar es la de incrementar el número de bosones en la teoría, lo que refuerza la transición de fase electrodébil. Sin embargo, nuestra estrategia es diferente y demostramos que si se cambia la historia térmica del Universo con la inclusión de un periodo de dominación de materia no relativista, el límite de esfalerón puede ser más laxo y la ventana para la generación de la asimetría bariónica abierta de nuevo.

Finalmente, un estadío intermedio de dominación de materia no relativista también puede tener consecuencias en la evolución de perturbaciones primordiales. Como es bien conocido, las perturbaciones dentro del horizonte crecen mucho más durante un periodo dominado por materia no relativista que durante radiación. En el escenario convencional, este crecimiento durante la dominación de la materia oscura condujo a la formación de los objetos que vemos. Pero esto crea otra cuestión: ¿Qué les ocurre a las perturbaciones si también hubo un periodo de materia no relativista antes de BBN?. Nosotros consideramos este escenario, que puede ser descrito completamente por las características de la(s) partícula(s) que dominan el Universo durante dicha etapa. Las ecuaciones de movimiento de las perturbaciones durante dicho periodo fueron resueltas y evolucionadas y mostramos que este periodo de materia no relativista anterior a BBN puede dejar trazas en forma de estructuras y subestructuras que en principio se podrían detectar.

La estructura de la Tesis

Esta Tesis pretende explorar nuevos escenarios complementarios a los estándar que se han considerado hasta el momento en la literatura. Esta Tesis consiste entonces en una serie de algunos artículos publicados durante mi doctorado y que han sido particularmente seleccionados para seguir una linea de investigación y exposición coherente.

Las partes de la Tesis que siguen a este capítulo son las siguientes:

- Una parte donde se introducen y revisan los conceptos mínimos que pueden ser necesarios para el entendimiento de los artículos. Esta parte a su vez ha sido dividida en cuatro capítulos. En el primero se presentan las características principales del modelo cosmológico estándar y la historia térmica del Universo. A continuación, se explica con gran detalle la física inflacionaria. El siguiente capítulo está dedicado a la generación de la asimetría bariónica en el Modelo Estándar. Por último, se concluye con una revisión minuciosa del marco de formación de estructuras.

- En la siguiente parte de la Tesis se presentan como tal los artículos en los que he trabajado. En el primer capítulo, se explica el escenario de inflación y generación de asimetría bariónica debido a un condensado de neutrinos dextrógiros pesados. En el siguiente capítulo, se estudia la transición de fase electrodébil cuando se introduce un periodo intermedio de dominación de materia no relativista. Se concluye con el capítulo en donde se aplica este escenario al estudio del crecimiento y amplificación de perturbaciones primordiales.

- Concluyo la Tesis con una parte dedicada a resumir los resultados fundamentales y conclusiones que pueden ser extraídas de los artículos. Asimismo, se muestra la literatura usada en los artículos y en la parte introductoria de revisión de los conceptos teóricos.

Part II

THEORETHICAL BACKGROUND

1. STANDARD COSMOLOGICAL SCENARIO: HOT BIG BANG MODEL

The standard model of particle physics provides an elegant description of the electroweak and QCD interactions and reproduces with great success and precision the phenomenology behind. Moreover, such a scenario has been extensively proved experimentally and with the recent discovery of the Higgs boson, many theoretical issues seem to have been addressed.

With regard to cosmology and the behaviour of the Universe, it is likewise accepted that the Hot Big Bang scenario establishes the standard framework for the evolution of the Universe as a whole. This model relies on General Relativity Theory in addition to the so called Cosmological Principle: "The Universe looks the same whoever and wherever you are". Further than being an assumption, the cosmological principle has been confirmed observationally proving the homogeneity and isotropy of the Universe at large scales.

The ingredients of our standard knowledge of the Universe are based on the following experimental evidences:

- Isotropy and homogeneity, measured at the CMB
- Expansion of the Universe, accelerated nowadays from Supernova observations
- Abundances of light elements
- Formation of galaxies and substructures

The success of the Big Bang scenario is its ability to address and match the aforementioned evidences with remarkable accuracy, giving us a considerable confidence on it. Moreover, the resulting picture from such a scenario turns out to be astonishingly simple. The Universe takes off at very early times from a very hot and dense gas which cools down as the Universe expands. As a result, energy density decreases and temperature falls. Likewise, the Universe consists of several components for which their evolution with expansion is quite different, as we will see later. This leads to the domination of these different species along different epochs in the history of the Universe. At first stages, the energy density is dominated by radiation, which includes the relativistic particles, and at later stages it is dominated by an anti-gravitational source of energy, commonly known as Dark Energy due to its unknown nature. In between, the Universe underwent a period dominated by Dark Matter, a weakly interactive and highly massive type of matter. Moreover, this well-defined description of the history of the Universe has been confirmed observationally.

1 The Friedman Robertson Walker Universe

The Cosmological principle can be expressed mathematically by means of the FLRW metric

$$ds^2 = dt^2 - a(t)^2 \left(\frac{dr^2}{1 - kr^2} + r^2(d\theta^2 + \sin^2\theta d\varphi) \right) , \tag{1.1}$$

which describes the Universe in terms of the expansion rate $a(t)$ and its spatial curvature k.

Furthermore, the FLRW metric also constrains the form of the stress energy tensor $T_{\mu\nu}$ to be diagonal and with equal spatial components. The simple way to accomplish these features is to consider a perfect fluid characterised by the time-dependant pressure $P(t)$ and energy density $\rho(t)$

$$T^{\mu\nu} = (\rho + P)u^\mu u^\nu + P\eta^{\mu\nu} , \tag{1.2}$$

where $\eta^{\mu\nu}$ is the Minkowsky metric and u^μ is the 4-velocity.

Using the metric and the stress energy Tensor, one can fully derive the Einsteins Equation $R_{\mu\nu} - \frac{1}{2}g_{\mu\nu}R = T_{\mu\nu}$. The $0-0$ components is especially important since it gives the well-known Friedman equation, which reads

$$H^2 \equiv \left(\frac{\dot{a}}{a} \right) = \frac{8\pi G}{3} - \frac{k}{a^2}\rho . \tag{1.3}$$

H is usually named the Hubble parameter and it basically gives the expansion rate of the Universe, which also allows us to know the length of the Universe as H^{-1}, *i.e* the maximum distance that light has travelled since the Big Bang. It is also referred sometimes as the Hubble horizon and given in comoving units as $(aH)^{-1}$. We will see that this last definition is very useful when talking about inflation.

In addition, the dynamics of the different components of the Universe's fluid can be obtained my making use of the covariant continuity equation $\nabla_\mu T^{\mu\nu} = 0$. Focusing on the energy density part, it is easy to demonstrate that the evolution of ρ is given by

$$\begin{aligned} \dot{\rho}_i &= -3H(\rho_i + P_i) \\ &= -3H(1 + w_i)\rho_i \end{aligned} \tag{1.4}$$

where w_i, defined as $w_i = \frac{P_i}{\rho_i}$, characterises the equation of state of a perfect fluid and is assumed to be constant unless interactions between different components exist. Thus, the evolution in time for a given i-fluid component is $\rho_i(t) \propto a^{-3(1+w_i)}$, which is different for each component. In particular

$$\begin{array}{ccll} w = 1/3 & \text{“Radiation”} & \Rightarrow & \rho_i(t) \propto a^{-4} \\ w = 0 & \text{“Matter”} & \Rightarrow & \rho_i(t) \propto a^{-3} \\ w = -1 & \text{“Cosmological Constant”} & \Rightarrow & \rho_i(t) \propto \text{constant} \end{array}$$

As a consequence, due to the different dilution of the fluid components, it is clear that the Universe went through different domination regimes throughout its thermal history.

2 *Brief description of the standard thermal history*

The departing point in the standard history of the Universe in the Hot Big Bang scenario begins somewhat earlier than BBN at a temperature of 1 MeV. The Universe composition and evolution in terms of cosmology before BBN is somewhat model dependent in relation to particle physics, which may lead to different scenarios. Moreover, BBN supposes the earliest experimental test in Cosmology.

- $T_{BBN} > T > T_{eq}$
  In this period of time, the Universe was dominated by relativistic matter ρ_R. The energy density in this case follows a blackbody distribution and is written as

$$\rho_R = \frac{\pi^2}{30} g_*(T) T^4 \ , \tag{1.5}$$

  where g_* accounts for the number of degrees of freedom in relativistic bosons and fermions and reads

$$g_*(T) \equiv \sum_{\text{bosons}} g_*^i + \frac{7}{8} \sum_{\text{fermions}} g_*^i \ . \tag{1.6}$$

  Around T_{BBN}, neutrinos decouple from the plasma made up of protons, neutrons and photons. As a consequence, β-decay stops operating efficiently and this allows the neutrons to be stable and finally bind with protons in a complex way to form the first light elements such as Hydrogen and Helium.

- $T_{eq} > T > T_0$
  Due to the rapid dilution of radiation compared to that of non relativistic matter, both energy densities cross each other at a temperature given in terms of today parameters as

$$T_{eq} = T_0 \left(\frac{\Omega_{M,0}}{\Omega_{R,0}} \right) \ , \tag{1.7}$$

  where $T_0 \simeq 2.73K$, $\Omega_{M,0} \simeq 0.26$ and $\Omega_{R,0} \simeq 10^{-5}$ are the current temperature of the Universe, matter and radiation density parameters respectively. Taking these values into account, one finds that the Matter-radiation equality temperature is $T_{eq} \simeq$ 1eV. Once below this point, the total energy density of the Universe falls like a^{-3}, as it was shown before for matter domination.

  At a $T \simeq 0.3$eV, photons are no longer able to ionise the plasma and electrons can bind with nuclei to form the first atoms. Later on, a remarkable event takes place near a temperature $T_{\rm dec} \simeq$ 0.1 eV, when reactions of the form $\gamma + e^- \longrightarrow \gamma + e^-$ stop being effective and radiation is then decoupled from matter. At this stage, the Universe becomes ”visible” and provides us with the Cosmic Microwave Background map, the first landscape of the early Universe.

Finally, after this decoupling point, baryons, which are no longer tight to photons, begin to fall in the gravitational potential wells previously originated by Dark Matter, collapse and form the first bound objects.

- $T_0 > T$
  The energy density of the Universe keeps diluting as the third power of the scale factor until the current day, when a new type of energy, called “Dark energy“ Ω_{DE} with an almost constant energy value, dominates the plasma. This yields the observed acceleration of the expansion of the Universe.

2. INFLATION

1 Issues of the Standard Cosmology Model which motivates Inflation

1.1 The Horizon Problem

The fact that the light speed is finite implies the existence of horizons where physics process are causally connected. From the CMB experiments, we see that the sky temperature map is vastly uniform, which means that at some moment in the history of the Universe thermodynamic equilibrium was established. Nonetheless, in the standard picture it is easy to demonstrate that the size of the Universe is much bigger than the comoving distance travelled by photons by the decoupling time.

Therefore, points of the sky separated by an angle of more than about a degree, could have never been in contact to establish thermodynamic equilibrium. Consequently, the Hot Big Bang model is unable by itself to offer an explanation for the homogeneity in temperature seen accurately in different regions of the sky.

1.2 Flatness Problem

The Friedmann equation, which relates the Hubble parameter with the energy density of the different components of the Universe, may be written as follows

$$\Omega(t) - 1 = \frac{k}{a(t)^2 \; H(t)^2} \,. \tag{2.1}$$

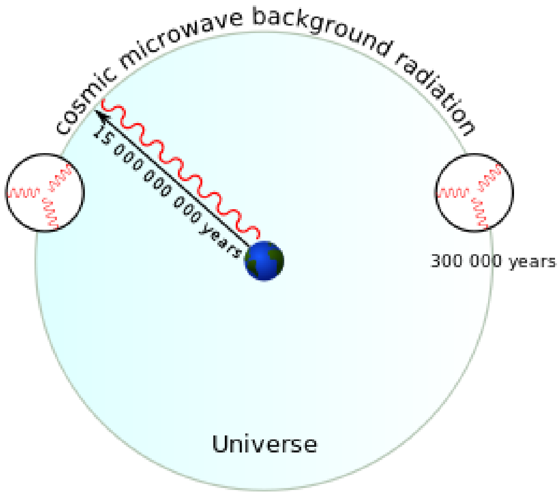

Fig. 2.1: Depiction of the horizon problem. The size of the Universe at the decoupling time is much bigger than the comoving distance travelled by photons from points well separated in the sky, which demonstrate that they were never in contact.

Current observations constrain $|\Omega(t_{\text{today}}) - 1|$ to be less than 0.01 [10], which implies that at earlier times, this difference was extremely smaller. For example, at the decoupling time

$$|\Omega(t_{\text{dec}}) - 1| \lesssim 10^{-16} , \tag{2.2}$$

and at the Planck epochs

$$|\Omega(t_{\text{Planck}}) - 1| \lesssim 10^{-60} . \tag{2.3}$$

Therefore, as we can see, $\Omega(t)$ was astonishingly close to the unity at early stages. The flatness problem states that such fine tuning in the initial conditions of the Universe seem exceptionally unlikely since any tiny deviation from this would lead to a different geometry of the Universe.

1.3 Unwanted particles

In the hot big bang picture, the Universe begins at a very high temperature, which makes that relics forbidden by observations may be produced kinematically and survive up to the present day, contributing to the density energy of the Universe. The Hot Big Bang model has no mechanism for getting rid of these relics produced at very early stages in the history of the Universe.

1.4 Small scale structure

We have mentioned that CMB experiments indicate that the Universe is quite homogeneous and smooth. However, this is only true on large scales since on smaller scales there exists a vast variety of structures: stars, galaxies, clusters. . . .

The Big Bang Model provides with a standard framework for the formation of structures from very small seeds. In addition, in order to explain the origin of such objects, one needs to assume the existence of initial inhomogeneities with a gaussian probability distribution and an almost scale-invariant power spectrum. However, the origin of these density inhomogeneities remains a mystery under the standard scenario and how such scale-invariant spectrum is produced.

1.5 Inflationary Universe as a solution

The definition of an inflationary Universe is simply: a Universe which undergoes an epoch where the scale factor grows exponentially. During such an epoch, the comoving Hubble length evolves as

$$\frac{d}{dt}\left(a(t)H(t)\right)^{-1} < 0 , \tag{2.4}$$

where $H(t)$ is the Hubble rate and $a(t)$ the scale factor. The above definition of inflation states that the observable Universe becomes smaller during the inflationary epoch.

Once given the definition of Inflation, one can see how all the issues of the Big Bang Theory are solved.

Firstly, the *horizon problem* is solved since the reduction of the Hubble length during Inflation allows our observable Universe to originate from a tiny region which was well inside the Hubble radius, *i.e*, casually connected process were possible, at that time.

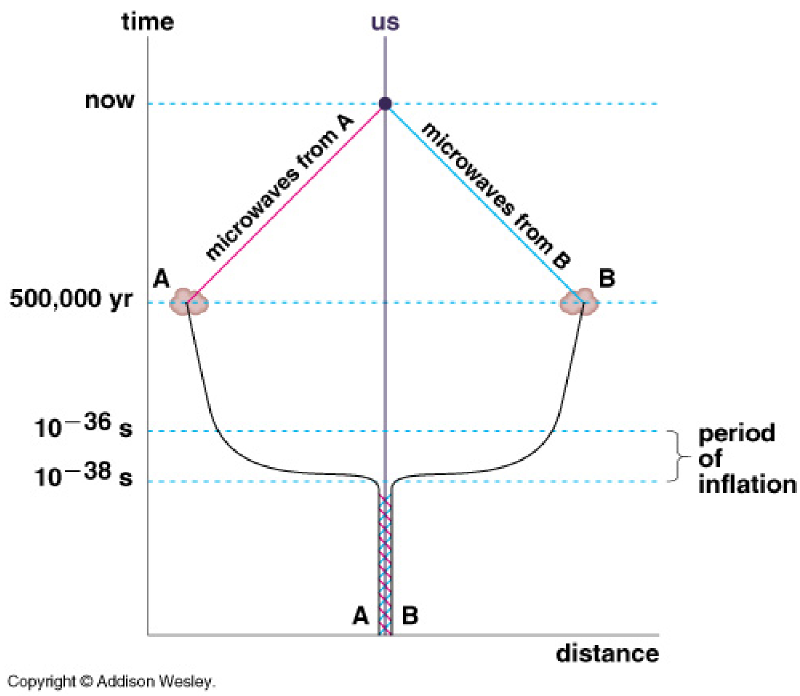

Fig. 2.2: Solution to the horizon problem by the introduction of exponential growth of the Universe

Secondly, during inflation the energy density of the Universe remains constant while the scale factor grows exponentially. Therefore, the curvature term $\frac{K}{a(t)^2H(t)^2}$ decreases to the tiny values given in the previous section. This is how the *Flatness problem* is solved.

On the other hand, relic abundances of *unwanted particles* produced before the inflationary epoch are diluted to a satisfactory level by the huge growth of the Universe during inflation.

Finally, the first seeds that originate all the rich *small scale structure* may arise due to the quantum fluctuations of a scalar field ϕ, which plays the role of the inflaton. As might be expected, the inflationary scenario might take place in more complex scenarios not restricted to a single field, but the origin of the first seeds is similar.

2 Inflation and scalar fields

In this section, I shall explain how the inflationary framework can be easily developed by introducing a scalar sector in the theory. For simplicity, I shall focus on the single field case, even though the general features of Inflation can be simply reproduced in more complex scenarios.

Inflation is driven by a scalar field ϕ, often called the inflaton. Such a field needs to possess certain features: negative pressure to accelerate the expansion of the Universe and live enough to solve all the problems and lead to the Big bang picture.

The first feature is easily achieved when one writes down the component of the density energy and pressure associated to the Inflaton:

$$\rho_\phi = \frac{1}{2}\dot{\phi}^2 + V(\phi) \tag{2.5}$$

$$P_\phi = \frac{1}{2}\dot{\phi}^2 - V(\phi) \; , \tag{2.6}$$

where $V(\phi)$ is the potential and a homogeneous scalar field $\phi \equiv \phi(t)$ is assumed. Yet there might exist inhomogeneities on this field but it can be demonstrated that the spatial gradients in ϕ are suppressed by the square of the scale factor, so such components are shifted away rapidly as long as they are relatively small. In case the spatial gradients are comparable with kinetic and potential terms, they may prevent Inflation from commencing.

The equations of motion within this scalar framework can be then easily derived and read as follows

$$\ddot{\phi} + 3H\dot{\phi} = -\frac{dV}{d\phi} \tag{2.7}$$

$$H^2 = \frac{8\pi}{3M_{Pl}^2}\rho_\phi \; . \tag{2.8}$$

Therefore, we can see that the way to satisfy the condition 2.4 for inflation to occur is ensured if

$$\dot{\phi}^2 < V(\phi) \; . \tag{2.9}$$

Such an state is easily fulfilled if the potential is flat enough. In general, this flatness condition needs not to be obeyed initially since a scalar field suitably displaced from the minimum of its potential can quickly approach such a condition.

Hence, a given model for inflation is supplied by the form of its potential and any theory intended to naturally embed an inflationary stage must be able to dynamically explain the existence of such a potential. Examples of different potentials that can be found in the literature are

$$V(\phi) = \sum \frac{c_i}{M^{(i-4)}}\phi^i \qquad \text{polynomial inflation} \tag{2.10}$$

$$V(\phi) = V_0 \left[1 \pm \cos\left(\frac{\phi}{\mu}\right)\right] \qquad \text{Natural inflation} \tag{2.11}$$

$$V(\phi) = V_0 \left[1 - \exp\left(-q\frac{\phi}{\mu}\right)\right] \qquad \text{“Exponential” inflation} \tag{2.12}$$

It will be showed that recent measurements at PLANCK [11] allow us to efficiently rule out most of the proposed models.

3 Scales and the amount of inflation

The amount of inflation is usually defined by the so-called ”number of e-foldings” N, which is the ratio between the scale factors at the end and beginning of inflation, taken in logarithm scale:

$$N \equiv \ln \frac{a_f}{a_0} \tag{2.13}$$

The number of efoldings determines the duration of the inflationary stage required to solve the aforementioned problems in the hot Big Bang picture.

As it was mentioned earlier, during inflation the Hubble length $1/(aH)$ decreases with time, a fact that exhibits that a certain scale may cross the horizon, being outside at the end of Inflation. A scale is to cross the horizon when $k = aH$. Therefore, a given scale left the horizon at some time during inflation to re-enter again somewhat later in the history of the Universe. This simple procedure allows us to estimate the number of efoldings a given scale spends outside the horizon until it reenters again

$$\frac{k}{(aH)_{\rm reent}} = \frac{a_* H_*}{(aH)_{\rm reent}} = \frac{a_*}{a_{\rm end}} \frac{a_{\rm end}}{a_{\rm reh}} \frac{a_{\rm reh}}{a_{\rm eq}} \frac{a_{\rm eq}}{a_{\rm reent}} \frac{H_*}{H_{\rm reent}} \,, \tag{2.14}$$

where the subscript "*" is the value when the scale first left the horizon, "end" the value at the end of inflation, "reh" at the reheating time, "eq" at the matter-radiation equality point (*i.e.* when both radiation and matter energy densities are equal) and "reent" indicates the value when the scale reentered the horizon again.

Then, the number of efoldings is

$$N_k = 62 - \log \frac{k}{(aH)_{\rm reent}} - \log \frac{10^{16}{\rm GeV}}{V_k^{1/4}} - \log \frac{V_k^{1/4}}{V_{\rm end}^{1/4}} - \frac{1}{3} \log \frac{V_{\rm end}^{1/4}}{\rho_{\rm reh}^{1/4}} \tag{2.15}$$

Therefore, a number of efoldings around 60 is sufficient to causally connect the today patches of the Universe and extremely flatten the geometry of space-time.

4 Quantum fluctuations and the origin of structures

Inflation provides an elegant explanation for the origin of the initial seeds that led to the formation of the known structures. In order to accomplish this, Inflation invokes quantum mechanics and perturbs the inflation field with a quantum inhomogeneous fluctuation

$$\phi(t, \mathbf{x}) = \phi_h(t) + \delta\phi(t, \mathbf{x}) \tag{2.16}$$

When performing this perturbation over the equation of motion for the inflaton field eq.(2.7), one may obtain the equation of motion for the perturbation in the Fourier space, which reads

$$(\delta\phi_k)^{\cdot\cdot} + 3H(\delta\phi_k)^{\cdot} + m^2(t)(\delta\phi_k) = 0 \,, \tag{2.17}$$

where $m^2(t)$ is the time dependent mass: $\left(\frac{k}{a}\right)^2 + V''$.

The above equation needs to be solved by making use of quantum mechanics. Since during Inflation $m^2 \ll H^2$, so the Inflaton mass can

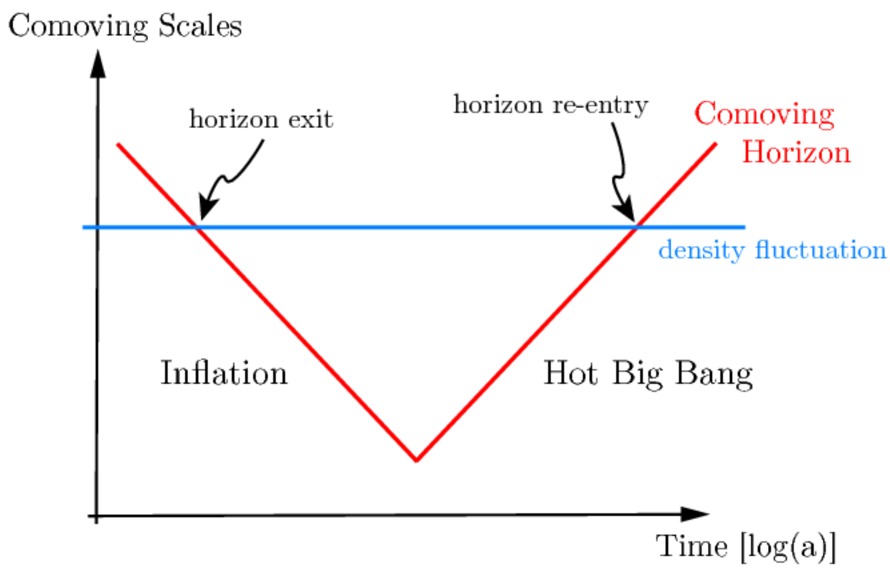

Fig. 2.3: Representative evolution of the different scales with time (scale factor). During Inflation, a given physical scale (in blue) leaves the horizon due to the decrease of the comoving Hubble length (in red).

be neglected, one can thus obtain, as in the case of the one dimensional harmonic oscillator, a non-zero variance for the Inflaton perturbations from the vacuum fluctuations as

$$\langle|\delta\phi_k|^2\rangle = \frac{H^2}{2k^3L^3}\left(1+\frac{k^2}{a^2H^2}\right) . \tag{2.18}$$

This result shows the way Inflation gives rise to a non-zero vacuum primordial fluctuation. As it was stated previously, the main requirement for an Inflationary stage is that the comoving Hubble length $(aH)^{-1}$ decreases in time. Therefore, a given scale k during inflation leaves the horizon at some point (see fig. 2.3) and the left hand side of equation 2.18 dominates, giving then

$$\langle|\delta\phi_k|^2\rangle = \left.\frac{H^2(t_*)}{2k^3L^3}\right|_{k=a(t_*)H(t_*)} , \tag{2.19}$$

where it is evaluated at the moment when the scale exits the horizon and L is a length factor. Nevertheless, during inflation H slightly changes, so the above equation can be simply evaluated at the time of horizon exit $k = aH$.

Once outside the horizon, causal physics stops being valid and different parts of the Universe evolve independently to each other as different Universes. Thus, the existence of the above non zero variance for the inflaton field on wavelengths exceeding the comoving Hubble Length makes the total Inflaton field (eq. 2.16) have slightly different values in these independent Universes. Hence, these different Universes may stop inflating at somewhat different times. This process converts the perturbations on the inflaton field into fluctuations on the energy density $\delta\rho$. However, the disadvantage of using $\delta\rho$ is that it is time-dependent on the scales in play. As a consequence, it is more useful to work with $\mathcal{R}$, which is the so-called spatial curvature generated by constant inflaton field hypersurfaces and which is constant on scales outside the horizon. Furthermore, such a parameter will be very useful when the structure formation picture is reviewed.

In order to calculate $\mathcal{R}$, one needs to make use of general relativity in a specific gauge, since on very large scales the geometrical changes induced by perturbations are important. It can be demonstrated that the curvature generated by the inflationary fluctuations, hereafter curvature perturbations, read in Fourier components as

$$\mathcal{R}_k = \frac{H}{\dot{\phi}} \delta\phi_k \; . \tag{2.20}$$

Hence, the primordial power spectrum of curvature perturbations $\mathcal{P}_{\mathcal{R}}(k)$ defined as the two point correlation function $\langle \mathcal{R}_k \mathcal{R}_{k'} \rangle$ is easily written as

$$\mathcal{P}_{\mathcal{R}}(k) = \left(\frac{H}{\dot{\phi}}\right)^2 \left(\frac{H^2}{2\pi}\right)\Bigg|_{k=a(t_*)H(t_*)} \simeq \frac{1}{12\pi^2 M_{Pl}^6} \frac{V^3}{V'^2}\Bigg|_{k=a(t_*)H(t_*)} \tag{2.21}$$

during Inflation.

Owing to the fact that different scales exit the Horizon at different times, the potential will have slightly different values, this leads to a small scale-dependence, parametrised trough the spectral index n_s and written as follows

$$n_s(k) - 1 \equiv \frac{d \log \mathcal{P}_\mathcal{R}}{d \log k} = 2 \left(M_{Pl}^2 \frac{V''}{V} \right) - 3 \left(M_{Pl} \frac{V'}{V} \right)^2 . \tag{2.22}$$

Likewise, a perturbation on the scalar field also induces the creation of tensor perturbations on the metric for which a power spectrum and spectral index can be defined as follows

$$\mathcal{P}_t = \frac{2}{3\pi M_{Pl}^4} V \bigg|_{k=a(t_*)H(t_*)} \tag{2.23}$$

$$n_t = 4 \left(M_{Pl} \frac{V'}{V} \right)^2 \tag{2.24}$$

When giving the experimental results, it is often used r instead of $P_t(k)$, which shows the ratio of the latter one with the scalar power spectrum

$$r \equiv \frac{\mathcal{P}_t}{\mathcal{P}_\mathcal{R}} = 8 M_{Pl}^2 \left(\frac{V'}{V} \right)^2 \tag{2.25}$$

As we can see, the parameters I have introduced are very important because they are deeply related to the features of the inflationary potential. Moreover, they are in principle measurable, so they allow us to probe experimentally the shape of the potential.

5 *Inflationary paradigm in April 2014*

In this section, I will try to briefly summarise the status of our current inflationary knowledge, which mostly come from measuring any anisotropic imprints on the CMB map.

Basically, in all cases what one seeks for is the reconstruction of the radiation multipole power spectrum for different angular scales, which raises information about the nature of primordial perturbations and the form of the inflationary potential. In particular, the temperature data recollected by PLANCK, which can be seen in figure 2.4 fits with the assumption of primordial perturbations being gaussian and adiabatic.

Furthermore, PLANCK data in addition to the data already recollected from other experiments allows us to find the most likely values for the scalar and tensor primordial power spectrum, which likewise can be used to reconstruct the inflationary potential. In figure 2.5, one can see the marginalised regions for n_s and r from PLANCK and other data sets in comparison to the theoretical predictions of some inflationary models.

Finally, measurements on the CMB polarisation are expected to be another important source of information for Inflation, specially to determine the existence or not of primordial tensor modes. Recently, the south pole telescope BICEP2 has reported the detection of primordial B-modes which would constrain the tensor to scalar ratio to $r \sim 0.2$ [12], a measurement which strongly supports the Inflationary paradigm. It is expected that in a short time, PLANCK will release a new set of data including the CMB polarisation, which will confirm the BICEP2 important results.

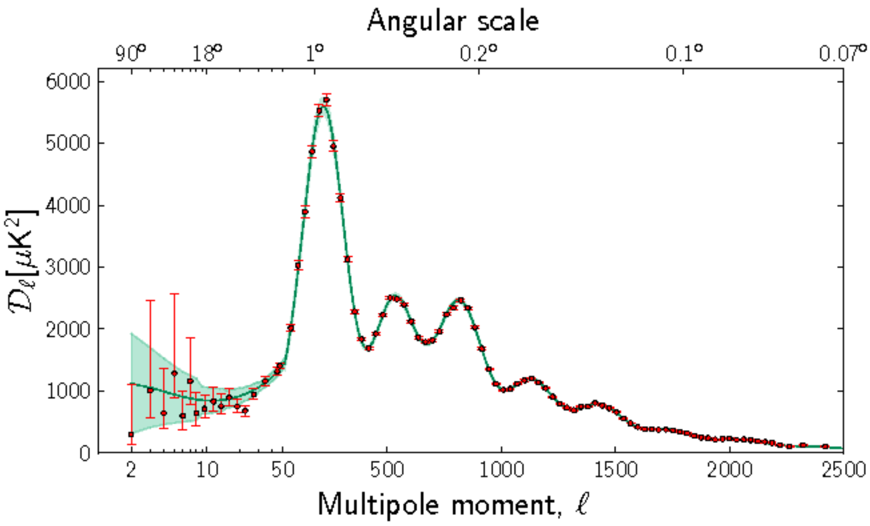

Fig. 2.4: Temperature power spectrum for different angular resolutions [13]

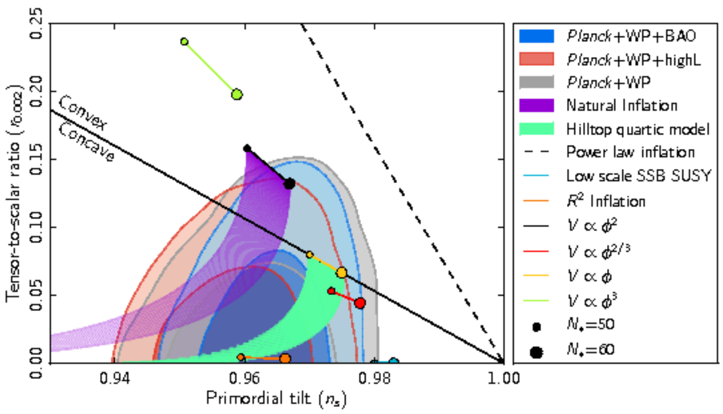

Fig. 2.5: Marginalised 68% and 95% confidence level regions for n_s and $r_{0.002}$ from PLANCK in combination with other data sets and compared to the theoretical predictions of selected inflationary models.

3. BARYOGENESIS

1 *Baryon asymmetry in the Universe*

In the previous chapter, the standard model of cosmology was introduced. As it was also outlined, it exhibits certain mysteries that such a model is not capable of addressing. One of these mysteries was associated with the number of baryons and antibaryons today. This quantity has been measured independently from the primordial nucleosynthesis abundances of light elements and WMAP results

$$\eta \equiv \frac{n_B - n_{\bar{B}}}{n_\gamma} \simeq 10^{-10} \tag{3.1}$$

2 *Sakharov's conditions*

In 1967 [14], Sakharov considered the question of how a non-zero baryon number could be originated from an initially symmetric state. Yet his original proposal for baryogenesis was long demonstrated to fail due to the proton instability, the conditions used are widely accepted to be sufficient (**but not necessary**) for any specific model of baryogenesis to give rise to a baryon asymmetry. Such conditions read as follows

- *B violation.*

  It is clear that as a minimum starting point there must exist a process $X \longrightarrow Y + B$ such that violates baryon number.

- *C and CP violation.*

  If C is conserved, the former reaction has the same width as its

C-conjugated reaction $\Gamma(X \longrightarrow Y+B)$=$\Gamma(\bar{X} \longrightarrow \bar{Y}+\bar{B})$. So, C should be violated otherwise both baryon numbers would get balanced over a long period of time.

Likewise, did we consider the B violating reaction which creates left handed particles $X \longrightarrow q_L q_L$, then right handed conjugate reaction would proceed at the same rate if CP were conserved and therefore

$$\Gamma(X \longrightarrow q_L q_L)+\Gamma(X \longrightarrow q_R q_R) = \Gamma(\bar{X} \longrightarrow \bar{q}_L \bar{q}_L)+\Gamma(\bar{X} \longrightarrow \bar{q}_R \bar{q}_R)$$

would still preserve baryon number. So, CP should be violated too along with C violation

- *Departure from thermal equilibrium.*

  If the Universe is in local thermal equilibrium, then the Boltzmann distribution for both matter and anti-matter with negligible chemical potentials should be the same since they have the same mass. Additionally, if a process that creates a net baryon number is out of thermal equilibrium, then the reversal process, *i.e.* the erasing net baryon number, is prohibited.

3 Baryogenesis in the SM

In the last section, I introduced the conditions needed to generate dynamically a baryon asymmetry beginning from a symmetric state. Therefore, it is natural to wonder whether such conditions are satisfied in the Standard Model under the current experimental constraints on the free parameters, specially, on the spectrum of masses and mixings.

$B+L$ *Violation*

Firstly, due to the chiral nature of the electroweak interactions, Baryon and Lepton number are highly violated in the SM when the Universe is in a thermal bath.

It can be demonstrated that the divergence of Baryon and Lepton currents in the SM are as follows

$$\partial^\mu J^B_\mu = \partial^\mu J^B_\mu = \frac{n_f}{32\pi^2}\left(-g^2 W^a_{\mu\nu}\tilde{W}^{a\mu\nu} + g'^2 B_{\mu\nu}\tilde{B}^{\mu\nu}\right) , \quad (3.2)$$

where n_f is the number of generations and $W^a_{\mu\nu}$ and $\tilde{B}^{\mu\nu}$ are the $SU(2)_L$ and $U(1)_Y$ field strength tensors respectively. Therefore, from these expressions one can see that $\partial^\mu(J^B_\mu - J^L_\mu) = 0$ and $\partial^\mu(J^B_\mu + J^L_\mu) \neq 0$, **B-L** is conserved in the SM whereas **B+L** is violated.

Integrating over space-time the above expressions for the divergence of the Lepton and baryon currents, it can be demonstrated that the difference in baryon and lepton number is

$$\Delta B = B(t_f) - B(t_i) = n_f \Delta N_{CS} \quad (3.3)$$
$$\Delta L = L(t_f) - L(t_i) = n_f \Delta N_{CS} , \quad (3.4)$$

where $N_{CS}(t) = \frac{g^3}{96\pi^2}\int d^3x \epsilon_{ijk}\epsilon^{abc}W^{ai}W^{bj}W^{cik}$ is the so-called Chern-Simons number with W^{ai} being a non-abelian gauge field. This number assigns a topological "charge" to the gauge fields associated with the ground state of our theory, with a rich structure due to its non-abelian nature.Consequently, transitions between different ground states which are separated by a certain potential barrier (see fig 3.1) can induce a baryon and lepton number difference of at least 3 units.

Such transitions can occur at two levels: "quantumly" and "classically". From a quantum point of view, they correspond to transitions through a tunnelling process and their rate can be demonstrated to be

$$\Gamma \sim e^{-\frac{4\pi}{\alpha_w}} \sim \mathcal{O}(10^{-165}) . \quad (3.5)$$

So we see that such a rate is extremely small, so **B+L** violation is negligible in the Standard Model at zero temperature. However, this violation does occur since the Standard Model is coupled to a thermal bath of temperature T and thermal fluctuations can induce transitions over the barrier depicted in figure between different ground states. This would correspond to the "classical level".

The finite-temperature transition between different vacua is determined by the so-called "sphaleron" configuration, which is an unstable

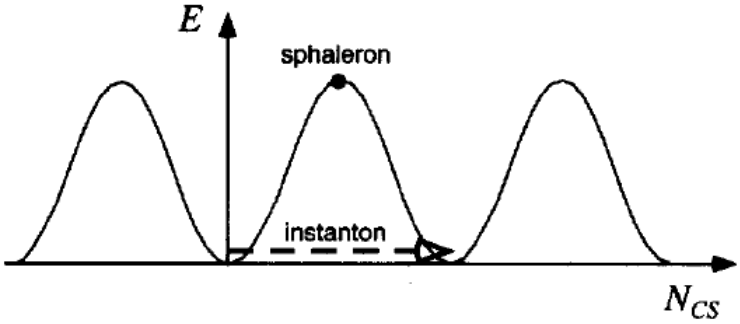

Fig. 3.1: Space configuration of the vacuum structure which shows the two possible paths from different vacua

solution in the gauge-Higgs space related to the minimum energy E_{sph} required to pass from one ground state to another violating B and L by equation 3.3 and 3.4. This energy can be written as follows

$$E_{sph}(T) \simeq \frac{8\pi}{g}\phi(T) \; , \tag{3.6}$$

where $\phi(T)$ is the Higgs vacuum expectation value that at zero temperature $\phi(0) \equiv v \simeq 246$ GeV.

As one can imagine from figure 3.1, when the energy is high compared to the barrier, transitions between different vacua can easily take place, since it would be as though there were no barrier whatsoever, *i.e.* $\phi(T) \simeq 0$, which corresponds precisely to $T \gtrsim 100GeV$ where the electroweak gauge symmetry is unbroken. At this stage, the transition rate is given naively by

$$\widetilde{\Gamma}_{\rm sph}(t) \propto T \; . \tag{3.7}$$

On the other hand, when the energy is smaller than the energy barrier E_{sph}, the transitions are suppressed by a Boltzmann factor $\exp(-E_{sph}(T)/T)$. For this last case, the transition rate takes the following naive form

$$\widetilde{\Gamma}_{\rm sph}(t) \propto \frac{\phi^7}{T^6} e^{-\frac{E_{\rm sph}}{T}} \; . \tag{3.8}$$

Since in this case $\phi(T) \neq 0$, such a transition rate is characteristic of the electroweak broken phase.

C and CP Violation

Due to the number of generations and the fact that particles have non degenerate masses in the SM, it is not possible to eliminate all the complex phases by redefinition of the fields in the quark sector. When performing such redefinitions over the fields, one finds the following charge current operator

$$\mathcal{L}_{cc} = -\frac{g}{\sqrt{2}}\bar{Q}\gamma^{\mu}V_{CKM}P_L d W_{\mu}^{+} + h.c. \ , \tag{3.9}$$

where $Q = (u_L, c_L, t_L)^T$ is the Quark $SU(2)_L$ doublet field, $d = (d_R, s_R, b_R)^T$ one of the two $SU(2)_L$ singlets field, $P_L = \frac{1-\gamma_5}{2}$, W_{μ}^{+} denotes the positive charge W gauge boson field and V_{CKM} is the so-called Cabbibo-Kobayashi-Maskawa mixing matrix, which is commonly written as follows

$$\begin{pmatrix} c_{12}c_{13} & s_{12}c_{13} & s_{13}e^{-i\delta_{CKM}} \\ -s_{12}c_{23} - c_{12}s_{23}s_{13}e^{i\delta_{CKM}} & c_{12}s_{23} - s_{12}s_{23}s_{13}e^{i\delta_{CKM}} & s_{23}c_{13} \\ s_{12}s_{23} - c_{12}c_{23}s_{13}e^{i\delta_{CKM}} & -c_{12}s_{23} - s_{12}c_{23}s_{13}e^{i\delta_{CKM}} & c_{23}c_{13} \end{pmatrix} , \tag{3.10}$$

where δ_{CKM} is the responsible for the CP violation.

Departure from thermal equilibrium

It was shown before that baryon number violation is mainly produced thermally through transitions between the different vacua of our theory. Such transitions violate baryon number by at least 3 units, which may be represented as a process that creates/destroys 9 left-handed quarks and 3 left-handed leptons (neutrinos for example). Such a process is called sphaleron process and is sketched in figure 3.2. We now need to wonder whether such process proceeds out of equilibrium as being required from the last Sakharov's conditions.

By comparing Γ_{sph} from eq. (3.7), *i.e.* in the Higgs symmetric phase, with the expansion rate H assuming a standard radiation dominated

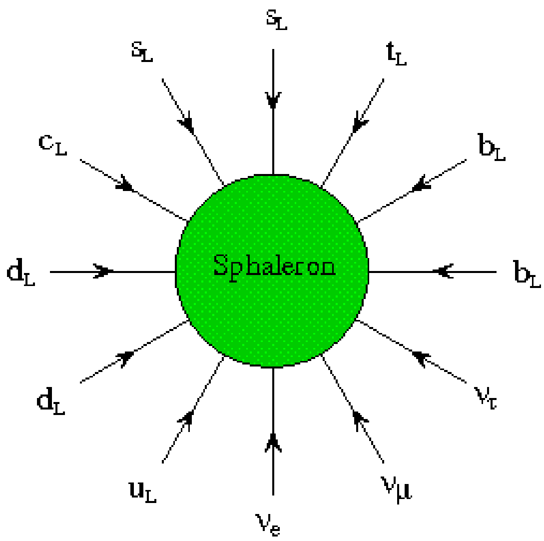

Fig. 3.2: Representation of a sphaleron process

Universe, one finds that when the temperature of the Universe is

$$10^{12}\ \text{GeV} \gtrsim T \gtrsim 100\ \text{GeV}\ , \tag{3.11}$$

the sphaleron processes are in thermal equilibrium and therefore they don't satisfy the third Sakharov condition and cannot be thus the source for the creation of baryon asymmetry.

It is now natural to turn our attention to the situation of thermal equilibrium in the non symmetric phase. As can be seen in the equation 3.8, in the Higgs broken phase $\phi(T) \neq 0$, so the exponential factor may slow down the transition rate sufficiently compared to the expansion rate to provide for the departure of thermal equilibrium, making viable the preservation of any baryon asymmetry created. On the other hand, such a suppression can also make the baryon and lepton violating processes effectively zero once the energy deeply drops, so baryogenesis must take place near the critical point that separates both phases. Thus, putting both requirements in turn, one finds that

$$\phi(T_c) > T_c\ , \tag{3.12}$$

where T_c is the critical temperature at which both phases coexist. The above bound, usually named as *sphaleron bound* in the literature, leads to two important conclusions:

- The electroweak phase transition needs to be first-order.
- The strength of the transition must be strongly enough to satisfy the sphaleron bound.

Finally, the remaining question to know is whether or not the above statements are satisfied within the Standard Model. In order to do so, one needs the effective Higgs potential that allows us to calculate the temperature dependent Higgs vacuum in terms of the parameters of our theory. In the one-loop approximation, such a potential for high temperatures can be written as follows

$$V(\phi, T) \approx \frac{M(T)^2}{2} - ET\phi^3 + \frac{\lambda_T}{4}\phi^4 \,, \tag{3.13}$$

where $M(T), B$ and λ_T are the temperature dependent effective mass, cubic term and quartic coupling respectively, which are all given at the one-loop ring improved level by

$$\begin{aligned} M(T) &= \sqrt{A(T^2 - T_0^2)} \,, \\ A &= \frac{2m_W^2 + m_Z^2 + 2m_t^2}{4v^2} + \frac{1}{2}\lambda_T \,, \\ E &= \frac{2}{3}\left(\frac{1}{2\pi}\frac{2m_W^3 + m_Z^3}{v^3} + \frac{1}{4\pi}\left(3 + 3^{\frac{3}{2}}\right)\lambda_T^{\frac{3}{2}}\right) , \\ \lambda_T &= \frac{m_H^2}{2v^2} - \frac{3}{16\pi^2 v^4}\left(2m_W^4 \ln\frac{m_W^2}{a_B T^2} + m_Z^4 \ln\frac{m_Z^2}{a_B T^2} - 4m_t^4 \ln\frac{m_t^2}{a_F T^2}\right) , \\ T_0^2 &= \frac{m_H^2 + 8\beta v^2}{2A} \,, \qquad \beta = \frac{3}{64\pi^2 v^4}\left(4m_t^4 - 2m_W^4 - m_Z^4\right) . \end{aligned} \tag{3.14}$$

with T_0 as the temperature at which the phase transition ends, $v = 246$ GeV as the usual Higgs vacuum expectation value at zero temperature, $a_B = (4\pi)^2 \mathrm{e}^{-2\gamma_E} \simeq 50$, $a_F = (\pi)^2 \mathrm{e}^{-2\gamma_E} \simeq 3.1$, and γ_E as the Euler's constant.

By making use of the above potential, one can see that in order for the phase transition to be strong enough to satisfy the sphaleron bound, the Higgs mass should be smaller than 50 GeV. Given the recent measurement of the Higgs mass at the LHC, such bound is far from being fulfilled and this is one of the reasons why the Standard Model fails to produce the right amount of baryon asymmetry.

4. STRUCTURE FORMATION

In a previous chapter, it was shown how the Inflationary scheme is capable of giving rise to the primordial seeds that led to the formation of the observable structures long afterwards. In the halfway, a complex process took place where initial perturbations grew in different stages until what we know nowadays. Such a topic can be found in the literature in many different ways. In most of them, Newtonian perturbation theory is first used specially when introducing the qualitative features that perturbations exhibit. However, such a framework is not general at all since it fails for scales outside the horizon. As a consequence, I will begin by revisiting the linear behaviour of scalar perturbations (only adiabatic) within the General Relativity framework, an approach to studying perturbation theory that has the benefit of being valid for any scale, providing us with a whole picture of the evolution of perturbations throughout the thermal history of the Universe. These perturbations in play are latterly responsible for the formation of the observable structures.

1 *Scalar perturbations and General Relativity*

Scalar perturbations refer to any perturbation performed over a scalar function F. Such quantity is usually given in terms of their Fourier components F_k

$$F(\mathbf{x},t) = \frac{1}{(2\pi)^{3/2}} \int F_k(t) e^{i\frac{\mathbf{k}\cdot\mathbf{x}}{a}} \,, \tag{4.1}$$

where k is the wavenumber and has units of $(\text{dimension})^{-1}$. This is the reason I will refer to "scale" when talking about a certain wavenumber.

In General Relativity, perturbations are defined assuming a certain geometry characterised by a metric. This is truly important when studying perturbations that are outside the horizon, *i.e.* $k < (aH)$, since at this level physics is very sensitive to the geometry of space-time. Thus, in addition to energy density, pressure and velocity, a given background metric tensor $g_{\mu\nu}$ is perturbed

$$g_{\mu\nu} \longrightarrow g_{\mu\nu} = g^0_{\mu\nu} + \delta g_{\mu\nu} \,, \tag{4.2}$$

where $g^0_{\mu\nu}$ is the background metric and chosen to be that of the Standard Cosmological Model, *i.e.* the FLRW metric. Furthermore, it can be demonstrated that this metric perturbation $\delta g_{\mu\nu}$ has 6 physical degrees of freedom. Two of these degrees of freedom couple to the density, pressure and the irrotational part of the velocity and are called *scalar perturbations.* As the reader can infer, I will just focus upon this kind of perturbations. The remaining degrees of freedom, that are beyond the scope of this chapter, are the *vector perturbations*, which couple to the sinusoidal velocity perturbation, and the *tensor perturbations*, which are responsible for the gravitational waves.

On the other hand, the definition of the metric is rather arbitrary due to the coordinate freedom of general relativity, which comes from the fact that once perturbed the Universe, is no longer homogeneous and therefore we do not have unique choice for the coordinate system. This is usually called *gauge freedom* and a particular choice of the metric is referred as a particular *gauge.* The gauge I will focus on is the so-called *Newtonian or longitudinal* gauge and is defined by the following form of the metric

$$\begin{aligned} ds^2 =& -(1+2\Phi)dt^2 + a(t)^2(1+2\Psi)d\mathbf{x}\cdot d\mathbf{x} \\ =& a(\tau)^2\left[-(1+2\Phi)d\tau^2 + (1+2\Psi)d\mathbf{x}\cdot d\mathbf{x}\right] \,, \end{aligned} \tag{4.3}$$

where in the last equation the conformal time $dt = a(\tau)d\tau$ has been introduced for future convenience.

The above particular choice of the metric is specially useful since the metric perturbation Φ becomes equal to the usual gravitational potential in the Newtonian limit. Furthermore, it turns out that in absence of anisotropies, which is true in perfect fluid case, we have

$$\Psi = -\Phi \,. \tag{4.4}$$

Working out the Einstein tensor $G_{\mu\nu}$ in the Newtonian gauge, it can be demonstrated that the most relevant Einstein equations are

$$k^2\Phi + 3aH\dot{\Phi} + 3a^2H^2\Phi = -\,4\pi G a^2 \sum_i \rho_i\delta_i \tag{4.5}$$

$$\Phi' + \mathcal{H}\Phi = -\,4\pi G a^2 \sum_i (\rho_i + P_i)v_i \tag{4.6}$$

$$\Phi'' + 3\mathcal{H}\Phi' + (2\mathcal{H}' + \mathcal{H})\Phi = 4\pi G a^2 \sum_i u_{s,i}^2 \rho_i\delta_i \tag{4.7}$$

where the sums are performed over all the i-components of the fluid,$' \equiv \frac{d}{d\tau}$, $\mathcal{H} \equiv \frac{a'}{a} = \dot{a}$, $\delta_i = \frac{\delta\rho_i}{\rho_i}$, $u_{s,i}$ is a characteristic sound speed for each fluid. Furthermore, making use of the above equations, one may also obtain a very convenient equation for the gravitational potential, which reads

$$\Phi'' + 3\mathcal{H}(1 + u_{s,i}^2)\Phi' + \left[2\mathcal{H}' - \mathcal{H}^2(1 - 3u_{s,i}^2)\right]\Phi + u_{s,i}^2 k\Phi = 0\,. \tag{4.8}$$

Likewise, from the perturbed relativistic energy conservation equation $\nabla_\mu T^{\mu\nu} = 0$, one may also obtain the following relevant equations for studying scalar perturbations

$$\delta_i + 3\mathcal{H}(u_{s,i}^2 - w_i)\delta_i - (1 + w_i)k^2 v_i = 3(1 + w_i)\Phi' \tag{4.9}$$

$$[(1 + w_i)v_i]' + \mathcal{H}(1 - 3w_i)(1 + w_i)v_i + u_{s,i}^2\delta_i = -(1 + w_i)\Phi \tag{4.10}$$

To conclude, the given set of differential equations determine the behaviour of density perturbations and the gravitational potential. In particular, one may study the relativistic matter perturbations during a Radiation dominated Universe (Non relativistic Matter perturbations will be thoroughly examined later due to their leading role in the formation of structures).

Within a Radiation dominated Universe, $\rho_{total} \approx \rho_{rad}$ and $\mathcal{H} = \frac{1}{\tau}$. Perturbations of relativistic matter δ_{rad} with $w = c_s^2 = 1/3$ source the equation for the gravitational potential, which in this case can be written as

$$\Phi'' + \frac{4}{\tau}\Phi' + \frac{k^2}{3}\Phi = 0\,, \tag{4.11}$$

whose solution in conformal time is rather simple and reads

$$\Phi(\tau) = -9\frac{\Phi_0}{(k\tau)^2}\left[\cos\left(\frac{k\tau}{\sqrt{3}}\right) - \frac{\sqrt{3}}{k\tau}\sin\left(\frac{k\tau}{\sqrt{3}}\right)\right]. \tag{4.12}$$

On the other hand, the equations for δ_{rad} and the velocity read

$$\delta'_{rad} - \frac{4}{3}k^2 v_{rad} = 4\Phi' \tag{4.13}$$

$$v'_{rad} + \frac{1}{4}\delta_{rad} = -\Phi \tag{4.14}$$

The solution for the above variables is

$$\delta_{rad} = \frac{6\Phi_0}{(k\tau)^3}\left[2\sqrt{3}((k\tau)^2 - 3)\sin\left(\frac{k\tau}{\sqrt{3}}\right) - k\tau((k\tau)^2 - 6)\cos\left(\frac{k\tau}{\sqrt{3}}\right)\right] \tag{4.15}$$

$$v_{rad} = -\frac{3\sqrt{3}k\Phi_0}{2(k\tau)^2}\left[2\sqrt{3}(k\tau)\cos\left(\frac{k\eta}{\sqrt{3}}\right) + ((k\tau)^2 - 6)\sin\left(\frac{k\tau}{\sqrt{3}}\right)\right] \tag{4.16}$$

2 *Adiabatic perturbations and their evolution in the superhorizon regime*

It can be seen that primordial perturbations produced in the simplest inflationary models are mainly adiabatic, which means that fluctuations only take place on the total energy density, making then temperature look inhomogeneously in space as CMB experiments state. Furthermore, the concept of adiabaticity implies no heat flow, so that energy cannot be transferred between the fluid components at the background level. Therefore, the changes in the energy density of every component must be the same, which can be expressed by making use of the energy continuity equation as follows

$$\frac{\delta_i}{1+w_i} = \frac{\delta_j}{1+w_j} \ . \tag{4.17}$$

Thus the relation between matter and radiation perturbations is

$$\delta_{matter} = \frac{4}{3}\delta_{rad} \ . \tag{4.18}$$

It is important to note that such adiabatic relations hold at any epoch as long as a certain mode is superhorizon.The reason is that in the subhorizon regime, causal physics works effectively and thus the components of the cosmic fluid may evolve in a different way throughout the thermal history of the Universe.

On the other hand, the fact that the evolution of the adiabatic energy density perturbations is the same in the superhorizon regime implies that the k-dependant curvature perturbation $\mathcal{R}_k$ remains constant. Such a quantity turns out to be extremely valuable since it is responsible for connecting primordial perturbations generated out of horizon during inflation to horizon entry at later epochs. It can be then demonstrated that in the Newtonian gauge, the relation between the curvature perturbation and the gravitational potential outside the horizon, for a given epoch with a dominating fluid with an equation of state w, is written as

$$\Phi_k = -\frac{3+3w}{5+3w}\mathcal{R}_k \, , \tag{4.19}$$

which is valid until horizon entry. Therefore, the potential is constant in the superhorion regime, with values

$$\Phi_k = -\frac{4}{6}\mathcal{R} \quad \text{in RD} \tag{4.20}$$

$$\Phi_k = -\frac{3}{5}\mathcal{R} \quad \text{in MD} \, . \tag{4.21}$$

In addition, equation 4.5 reads now

$$k^2\Phi + 3a^2H^2\Phi = 4\pi G a^2 \delta_{total}\rho_{total} \, . \tag{4.22}$$

Hence, for modes outside the horizon $k << aH$, it is straightforward to see that

$$\delta_{rad} = -2\Phi \quad \text{in RD} \tag{4.23}$$

$$\delta_{matter} = -2\Phi \quad \text{in MD} \tag{4.24}$$

which in turn imply that adiabatic perturbations following the relations in eq. 4.17 are constant in the superhorizon regime.

As a last comment, primordial perturbations can also have a contribution from isocurvature modes, *i.e* modes that, unlike adiabatic perturbations where the changes are in the total energy density, have non vanishing density fluctuations such that the total $\mathcal{R}_k = 0$. Furthermore, both adiabatic and isocurvutare solutions are independent of each other and then, allow us in our linear theory to express any initial information in terms of both modes. However, the generation of primordial isocurvature modes rely on exotic mechanisms beyond the simplest, and yet most likely from observational experiments, inflationary models and consequently fall beyond the scope of our studied scenarios.

3 Dark Matter perturbations at Matter domination

Dark Matter perturbations δ_{DM} deserve a detailed analysis. They were primarily responsible for the creation of the gravitational potential wells that allowed the accretion of matter and led to the formation of the today observable structures.

In this chapter I shall focus upon large Dark Matter scales that entered the horizon when the Universe was dominated by non relativistic matter *i.e.*, the pressure and velocity of the main component of the cosmic fluid are both negligible. In equation 4.7, it is straightforward to demonstrate that in a matter dominated Universe $a \propto \tau$, $(2\mathcal{H}' + 2\mathcal{H}) = 0$, so the equation for the gravitational potential reads

$$\Phi'' + 3\mathcal{H}\Phi' = 0 \; . \tag{4.25}$$

The solution to this equation is $\Phi = \Phi_0$, *i.e.* the potential does not dependent on time, in contrast to the situation in the radiation domination regime, where the potential falls quickly with time (eq. 4.12).

In addition, at this stage the term $\rho_{dm}\delta_{dm}$ dominates the right side of equation (4.5). Therefore, the equation for δ_{dm} is simply given by

$$\delta_{dm} = -\frac{1}{4\pi G a^2 \rho_{dm}} \left(k^2 + 3a^2H^2\right) \Phi \; . \tag{4.26}$$

At the superhorizon regime $k \ll aH$ and the above equation states then, as expected, that Dark Matter density perturbations are time-independent with a value $-2\Phi_0$.

On the other hand, the most noticeable effect takes place in the sub-horizon regime, where the first term in eq.(4.26) dominates and consequently Dark Matter density perturbations grow in time

$$\delta_{dm} = \frac{1}{3}\frac{a(\tau)}{a_*}\delta_0 \qquad k \gg a\ H\ , \tag{4.27}$$

where a_* is defined at the time that a given scale enters the horizon $a(\tau_*)H(a_\tau *) = k$.

In general, in a matter dominated Universe, such a linear behaviour in time is characteristic of non relativistic matter perturbations that weakly interact with radiation. Strong interactions may prevent matter from growing, a feature exhibited by baryons after the matter-radiation equality point which will be briefly discussed later.

4 Dark Matter perturbations at Radiation domination

I shall now concentrate on the small scales that entered the horizon when the Universe was dominated by radiation. The equations of motion for these perturbations read from the last sections as follows

$$\delta'_{DM} - k^2\delta_{DM} = 3\Phi' \tag{4.28}$$

$$v'_{DM} + \frac{1}{\tau}v_{DM} = -\Phi \tag{4.29}$$

where the gravitational potential is sourced by radiation perturbations and takes the form given in equation eq. 4.12.

The solution for δ_{DM} after a bit of calculus is

$$\delta_{DM}(\tau) = \delta_{DM,0} + 3(\Phi(\tau) - \Phi_0) - k^2\int_0^\tau d\tau'\tau'\Phi(\tau')\log\frac{\tau}{\tau'}\ . \tag{4.30}$$

As it was previously shown, the gravitational potential rapidly decays with time within the radiation dominated regime. Consequently, the integral converges and the integration can be easily performed. Furthermore, invoking the adiabatic relations in the initial data, the final result for δ_{DM} is

$$\delta_{DM}(\tau) = -9\Phi_0\left(\log\frac{k\tau}{\sqrt{3}} + \mathcal{C} - \frac{1}{2}\right)\ , \tag{4.31}$$

where $\mathcal{C} = 0.577\ldots$ is the Euler constant. Therefore, Dark matter perturbations during a radiation dominated Universe grow logarithmically with the scale factor due to the presence of pressure from radiation. As it can be seen, such a growth is not as prominent as in the matter domination regime, where it was shown that $\delta_{DM} \propto a(t)$.

The logarithmic growth of Dark matter density perturbation generates an extra-contribution to the gravitational potential

$$\Phi_{DM} = -\frac{a^2(\tau)}{k^2} 4\pi G \rho_{DM} \delta_{DM} \; . \tag{4.32}$$

It turns out hence that at some point, because of the rapid dilution of radiation energy density, $\delta_{rad}\rho_{rad} < \delta_{dm}\rho_{dm}$ and therefore the gravitational potential evolves along with the Dark matter perturbations. In this case, the equations of motions read

$$\delta'_{DM} - k^2 v_{DM} = 3\Phi'_{DM} \tag{4.33}$$

$$v'_{DM} + \mathcal{H}\; v_{DM} = -\,\Phi_{DM} \tag{4.34}$$

Such a system of equations can be easily transform into a second order differential equation for δ_{DM}

$$\delta''_{DM} + \mathcal{H}\; \delta_{DM} - 4\pi G a^2 \rho_{dm} \delta_{DM} = 0 \; , \tag{4.35}$$

where equation (4.32) has been used.

If the following convenient variable $y = \frac{a}{a_{eq}}$ is introduced, the above equation is written as follows

$$y(y+1)\frac{d^2\delta_{DM}}{dy^2} + \left(1 + \frac{3}{2}y\right)\frac{d\delta_{DM}}{dy} - \frac{3}{2}\delta_{DM} = 0\,. \tag{4.36}$$

Such an equation is usually known in the literature as the Meszaros equation and governs the evolution of Dark matter perturbations *at any regime* once radiation perturbations become negligible.

The solution to this equation is

$$\delta_{DM}(y) = \delta^{(1)}_{DM}(y) + \delta^{(2)}_{DM}(y) \; , \tag{4.37}$$

where

$$\delta_{DM}^{(1)}(y) = C_1 \left(1 + \frac{3}{2}y\right) \tag{4.38}$$

$$\delta_{DM}^{(2)}(y) = C_2 \left(1 + \frac{3}{2}y\right) \left(\log \frac{\sqrt{1+y}-1}{\sqrt{1+y}+1} + 6\frac{\sqrt{1+y}}{2+3y}\right) . \tag{4.39}$$

At small y, the above equation takes the following simple form

$$\delta_{DM}(\tau) = C_1 + C_2 \log \frac{\tau}{\tau_{eq}} , \tag{4.40}$$

with

$$C_1 = -9\Phi_0 \log k\tau_{eq}, \quad C_2 = -9\Phi_0. \tag{4.41}$$

since this equation has to match equation (4.31) where radiation was the leading gravitational source.

Finally, at late times ($y \gg 1$)

$$\delta_{DM}(\tau) = C_1 \frac{3}{2} y . \tag{4.42}$$

which shows that it grows linearly with the scale factor as expected once deeply in the matter domination regime.

5 Matter Power Spectrum today

One of the most decisive keys of the structure formation picture presented so far is to seek for their imprints and compare them to observations. Such a task is done by the construction of the power spectrum of Matter today $P_M(k, a_0)$, which basically defines the correlation between different scales of matter perturbations (both dark matter and baryons). This quantity, following the standard way introduced when studied the primordial power spectrum from inflation, is defined by

$$\langle \delta_M(\mathbf{k}, a)\delta_M(\mathbf{k}', a)\rangle = k^3 \frac{\mathcal{P}_M(k, a)}{(4\pi)} \delta(\mathbf{k} + \mathbf{k}') = \frac{P_M(k, a)}{(2\pi)^3} \delta(\mathbf{k} + \mathbf{k}') , \tag{4.43}$$

where $k \equiv |\mathbf{k}|$ and the subscript M generally includes both dark matter and baryons information. However, I will only focus my attention on

Dark Matter information for sake of simplicity. The effects of including baryons in the Matter power spectrum will be briefly mentioned in the following section.

Typically, we are interested in presenting the above quantity in terms of the primordial power spectrum $\mathcal{P}_\mathcal{R}$ which raises information about the first gravitational wells that at some point started to accrete matter and form structures. As it was already pointed out, information about the primordial inhomogeneities is provided by Inflation. The structure formation picture simply gives the link between the Inflationary seeds and what we observe nowadays at different scales.

As I widely explained in the previous sections, the behaviour of scales changes dramatically depending upon the time they enter the horizon. Consequently, the shape of the power spectrum may be different for the scales in study. In the following, I will try to give a qualitative description of this in terms of the magnitude of the scales.

- $k < k_{today}$. These scales are still out of the horizon

$$\delta_k(t_0) = \delta_0 \Longrightarrow \mathcal{P}_M \sim \mathcal{P}_\mathcal{R} \tag{4.44}$$

- $k_{eq} \gg k > k_{today}$. These scales entered the horizon at matter domination

  As it was shown, a given scale k entering at this regime begins to grow linearly with the scale factor and such a growth is given by equation (4.26). Thus, the power spectrum can be straightforwardly written as

$$\mathcal{P}_M(k, a_0) = \frac{k^4}{(4\pi G)^2 a_0^4 \rho_{dm,0}^2}\left(\frac{4\pi\Phi^2}{k^3}\right) = \frac{4}{25}\frac{k^4}{a_0^4 H_0^4 \Omega_{dm}^2}\mathcal{P}_\mathcal{R} \ , \tag{4.45}$$

  where the relation between the potential and the curvature in matter domination $\Phi = \frac{3}{5}\mathcal{R}$ has been used. Therefore, as expected due to the linear growth of matter perturbations at this stage, the power spectrum for a nearly scale invariant primordial perturbations ($n_s \approx 1$), rises abruptly with scales in the form

$$\mathcal{P}_M(k, a_0) \propto k^4 \tag{4.46}$$

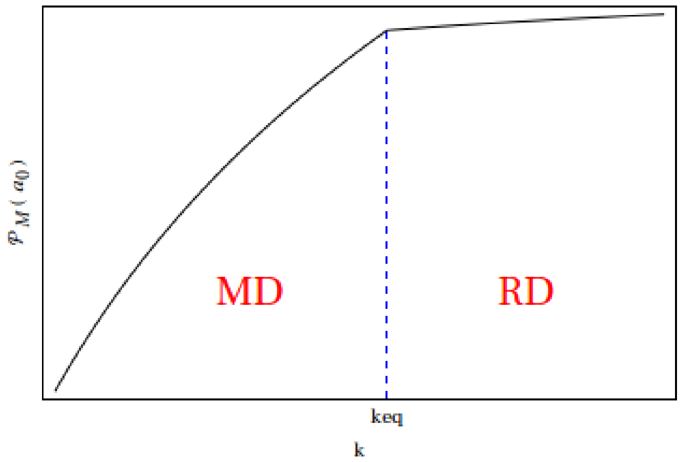

Fig. 4.1: Qualitative representation of the dimensionless power spectrum of Dark Matter perturbations $\mathcal{P}_M$ evaluated at a scale factor a_0 for scales that entered the horizon during the different eras of the thermal history of the Universe

- $k_{eq} \ll k$. These scales entered the horizon at Radiation domination

  In this regime, radiation is the main component of the Universe and the pressure associated prevents matter perturbations from growing, becoming almost flat during such a regime. The solution for these modes are given by the *Mezsaros* equation at late times and then the density contrast is mainly given by equation (4.42), so the power spectrum behaves as

$$\mathcal{P}_M(k, a_0) \propto \log \frac{k}{k_{eq}} , \tag{4.47}$$

  where one of the proportionality factors is $\left(\frac{k_{eq}}{a_0 H_0 \Omega_{dm}^{1/2}}\right)^4$ accounting for the linear growth that these modes experience once the Universe transits to the Matter domination epoch.

The behaviour of scales in both regimes is plotted in fig. 4.1, where the turnover around k_{eq} is smooth matching asymptotically both solutions.

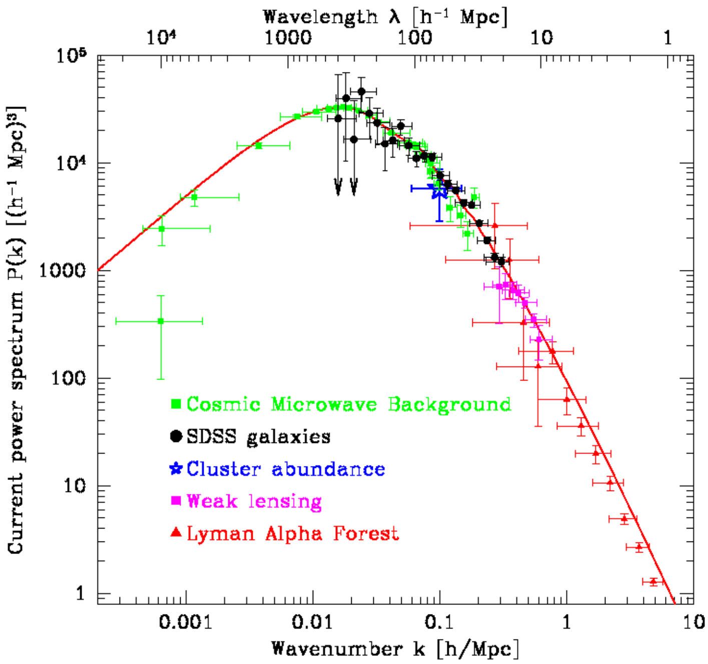

Fig. 4.2: Power Spectrum of Matter perturbations [15]

In the literature however, for observational purposes, it is more often used the dimensionful power spectrum $P_M(k, a_0)$, whose definition in terms of δ_M has been given above. The only change from $\mathcal{P}_M(k, a_0)$ comes into a suppression of k^3, which implies that

$$P_M(k, a_0) \propto k \qquad k \ll k_{eq} \tag{4.48}$$

$$P_M(k, a_0) \propto k^{-3} \log \frac{k}{k_{eq}} \qquad k \gg k_{eq}\ , \tag{4.49}$$

a profile exhibited in fig. 4.2, which is given in physical scales and where the turnover is located at $k_{eq} \sim 0.011(h\ \mathrm{Mpc}^{-1})$.

6 Last comments on structure formation

This last section is intended to make a few comments not included in previous sections as being considered far from the aim of this chapter but yet they deserve a brief mention.

- All this exposition of the standard structure formation scenario has been done in the linear regime, *i.e.* when $\delta\rho \ll \rho$. During this regime, one can work out the equations of motion for perturbations at leading order, which really simplifies the structure formation scheme, raises important information about it and still fit reasonably well with observations. The non-linear regime appears at really small scales that entered first in the horizon and had more time to grow. The computation of such perturbations during this regime needs to be performed using non-linear methods such as N-body simulations.

- In the qualitative calculation for the matter power spectrum provided in last section, only effects from Dark Matter have been considered. A more thorough calculation should include baryons as well, whose major effect is the decrease of the slope of the power spectrum and the introduction of oscillations at roughly scales $k > k_{eq}$.

- Likewise, a today matter power spectrum should also include the effects from Dark Energy, whose major effect is suppression at late times of the gravitational potential due to the stretching of space.

- To end with, it is important to point out that all the analysis performed in this chapter has been done assuming that Dark matter is non-relativistic when decoupled from the thermal bath, as observations seem to confirm. Furthermore, the fact about dealing with CDM implies the **bottom-up** picture of structure formation where the growth of structures proceeds hierarchically with smaller objects merging to form bigger structures.

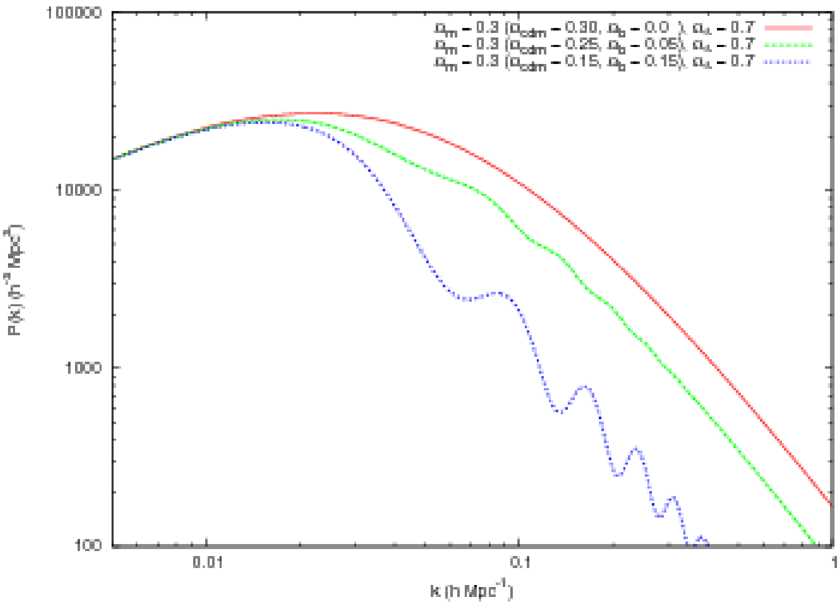

Fig. 4.3: Effects of baryons on the Matter Power Spectrum [16]

Part III

PUBLICATIONS

5. BARYOGENESIS FROM A RIGHT-HANDED NEUTRINO CONDENSATE

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abstract

We show that the baryon asymmetry of the Universe can be generated by a strongly coupled right handed neutrino condensate which also drives inflation. The resulting model has only a small number of parameters, which completely determine not only the baryon asymmetry of the Universe and the mass of the right handed neutrino but also the inflationary phase. This feature allows us to make predictions that will be tested by current and planned experiments. As compared to the usual approach our dynamical framework is both economical and predictive.

1 Introduction

There is nowadays an overall consensus that modern cosmology, based on the hot Big-Bang model and general relativity, constitutes a consistent theoretical framework which agrees *quantitatively* with data. It describes with amazing precision the evolution of the Universe from the first fraction of a second onwards. Nevertheless, such an impressive framework falls short of explaining the flatness and homogeneity of space, let alone the origin of matter and structures we observe in the universe today. As a result, no decent theory of the Universe lacks a judicious period of inflation, which wipes out the above mentioned problems.

However, despite its wide use, inflation is far from being a theory. Inflation is just a set of models of the very early universe which involve a period of exponential expansion, blowing up an extremely small region to one equivalent to the current horizon size in a fraction of a second. While the detailed particle physics mechanism responsible for inflation is not known, the basic picture makes a number of predictions that have been confirmed by observation. Inflation is thus now considered part of the standard hot Big Bang cosmology.

There are a bewildering variety of different models to realize inflation. In most of them however, inflation is parametrized through a single scalar field that fills space and which is assumed to have a potential energy. For a scalar field the total energy density and pressure are given by

$$\epsilon = 1/2 \left(\dot{\phi}^2 + V(\phi)\right) \tag{5.1}$$

$$p = 1/2 \left(\dot{\phi}^2 - V(\phi)\right) \tag{5.2}$$

If the field is changing slowly, so that the kinetic terms are much smaller than the potential ones, then we have $p \simeq -\epsilon$ and thus a component that can produce exponential expansion if it dominates the total energy density. Successful inflation thus requieres a phase in which the potential energy dominates the energy and pressure budget for a sufficiently long time.

Models of inflation differ in the assumed physical significance of the field ϕ, which is almost universally considered as a fundamental scalar. Although very popular, specially in particle physics where they plague most theories beyond the Standard Model, it is important to keep in mind that so far no fundamental scalar field has been observed. Thus, alternatives to a fundamental scalar have been looked for. Specially interesting paths were developed by technicolour [1,2], extended technicolour [3,4], walking technicolour [5] and topcolour [6]. Following this path, one of us [17] has recently pointed out the possible nature of the inflationary scalar field as a composite of massive right handed neutrinos.

At first, the existence of heavy right handed neutrinos, with trivial quantum numbers under the SM group, provides the simplest explanation for the origin of the neutrino mass (massless neutrinos go hand in hand with the absence of right-handed neutrinos). In order to make right handed neutrinos heavy, we have to allow them to develop Majorana masses, i.e. to give up the difference between neutrino matter and anti-matter.

Right handed neutrinos do not interact via electromagnetic, strong or weak interactions, instead they only mix with the light SM neutrinos (via the seesaw mechanism) in such a way that the observed mixture becomes massive. According to the simple seesaw model of mixing, the mass of the light neutrinos is of $\mathcal{O}\,(m_D^2/M_{RH})$, where M_{RH} is the mass of the heavy neutrino and m_D is a typical SM Dirac mass. The mere existence of Majorana fields, induces lepton number violation processes. This feature will play a fundamental role in our analisis.

As if the situation were not puzzling enough, it is remarkable the no observational presence of antimatter in the Universe. Several measurements coming from BBN, CMB and SNIa quantize this asymmetry by

$$\eta \equiv \frac{n_B - n_{\bar{B}}}{n_\gamma} \approx 10^{-10}, \tag{5.3}$$

where η denotes the asymmetry between baryons n_B and antibaryons $n_{\bar{B}}$, normalized to the number of photons n_γ. As usual, the market offers a wide array of mechanisms to address this quantity. Basically, they are based upon fullfilling the so-called Sakharov conditions, which are sufficient but not neccesary to generate dynamically this asymmetry. The different scenarios span from generation of the baryon asymmetry due to decaying GUT particles, baryogenesis produced by quarks reflections in front of Higgs bubbles during a first order electroweak phase transition or leptogenesis (for an excellent review see [18]). Along a completely different track, Cohen and Kaplan [19] demonstrated that the existence of spontaneous CPT violation in the theory by means of a derivative coupling of a scalar field to a baryon current permits the generation of the baryon asymmetry in equilibrium, without CP violation. In this work, we will go along with this path.

Our work is organized as follows: section 2 reviews how the addition of a four-fermion self-coupling of the right handed neutrino, if strong enough, triggers spontaneous breaking of the lepton number and produces a Majorana mass for the right handed neutrino. The cosmological implications of the effective potential (generated at one loop level) for the condensate are analyzed in section 3. We show in section 4 that, due to lepton number violation processes, one can produce a net lepton number density when the inflaton decays into ordinary matter. During the electroweak phase transition, such a net lepton asymmetry is converted into a baryon one via sphalerons proccesses. We discuss in section 5 the results obtained and conclude.

2 A right handed solution for the scalar field

In this section, we summarize the basic features of the model under which the calculation of the baryon asymmetry will be performed.

We would be interested in providing a dynamical origin to the scalar field, with a vacuum expectation value close to the energy scale of inflation. In order to do so, an effective four fermion self-coupling of the right handed neutrino field of strength G will be introduced by hand. This new interaction, should be strong enough to form a neutrino condensate that will trigger spontaneous symmetry breaking of lepton number and produce a Majorana mass for the right-handed neutrino. Below the cutoff scale Λ, the high frequency modes of the right handed neutrinos can be integrated, obtaining an effective theory of a Higgs-like composite field, which mimics the inflaton.

Such a four-fermion self interaction takes the form

$$G(\bar{\nu}_R^c \nu_R)(\bar{\nu}_R \nu_R^c), \tag{5.4}$$

where G is the dimensionful coupling constant, ν_R is the right handed neutrino and ν^c indicates charge conjugation. This is an effective interaction describing the physics below the cutoff Λ. There may be other higher dimension operators, but these will have subdominant effects at energies substantially below the cutoff scale.

In the limit of a large N_F, where N_F is the number of right handed neutrino flavours under the new interaction, there will be a solution to the gap equation for the dynamically induced right handed neutrino mass,

$$\begin{aligned} m_R &= -\frac{1}{2} G \langle \bar{\nu}_R \nu_R^c \rangle \\ &= -2GN_F \int \frac{d^4 l}{(2\pi)^4} (-1) \mathrm{Tr} \left(\frac{i}{\not{l} - m_R} \right). \end{aligned} \tag{5.5}$$

when

$$G\Lambda^2 \geq \frac{8\pi^2}{N_F}. \tag{5.6}$$

When this condition is satisfied, the theory predicts a scalar bound state with a mass of order m_R (to leading order in $1/N_F$). This is a standard result quoted for the Nambu-Jona-Lasinio model. It is important to stress that this bound state is a physical observable boson.

This physical particle is a bound state of $\bar{\nu}_R \nu_R^c$, arising by the attractive four-fermion interaction at the scale Λ of equation (5.4). This composite-boson $\Phi(x) = \rho(x) e^{i\frac{\phi(x)}{v}}$ is a complex field, with $\rho(x)$ its radial part, $\phi(x)$ the phase field and v an energy scale we will identify with a vacuum expectation value(vev). This parametrization shows that the right number of the degrees of freedom is kept after right handed neutrino condensation at scales below Λ.

In terms of the new particle, we can rewrite equation (5.4) as

$$g_o \left(\bar{\nu}_R^c \nu_R \Phi + \text{h.c.} \right) - m_0^2 \Phi^\dagger \Phi. \tag{5.7}$$

Notice that the new effective scalar field does not have a kinetic term, and it reproduces the four fermion vertex as an induced interaction when integrated out, with the identification

$$G = \frac{g_o^2}{m_o^2}. \tag{5.8}$$

To study the low-energy dynamics, we use the renormalization group to define effective low-energy couplings. This way, the running couplings at the scale μ are defined by integrating out all momentum-space degrees of freedom with momenta greater than μ. As we run down from the scale Λ downward in energy, all the possible couplings consistent with symmetries will be generated. However, it is expected that at scales below Λ, the theory can be parametrized by an effective Lagrangian which contains only "relevant" operators, with canonical mass dimension of four or less. In our case this means that the scalar field develops induced, fully gauged-invariant, kinetic terms and quartic term self-interactions from loop corrections, giving the renormalized lagrangian :

$$\mathcal{L} = \mathcal{L}_\Phi + \mathcal{L}_{SM} \tag{5.9}$$

with

$$\mathcal{L}_\Phi = Z\partial_\mu\Phi\partial^\mu\Phi^\dagger + g_o\left(\bar{\nu}_R^c\nu_R\Phi + \text{h.c.}\right) - m_\Phi^2\Phi^\dagger\Phi - \lambda_0\left(\Phi^\dagger\Phi\right)^2 \ , \tag{5.10}$$

where

$$Z = \frac{N_F\ g_0^2}{(4\pi)^2}\ \ln\left(\frac{\Lambda^2}{\mu^2}\right) \tag{5.11}$$

$$m_\Phi^2 = m_o^2 - \frac{2\ N_F\ g_o^2}{(4\pi)^2}\left(\Lambda^2 - \mu^2\right) \tag{5.12}$$

$$\lambda_o = \frac{2\ N_F\ g_o^4}{(4\pi)^2}\ln\left(\frac{\Lambda^2}{\mu^2}\right). \tag{5.13}$$

and $\mathcal{L}_{SM}$ the standard model (SM) Lagrangian which contains, among others, a Dirac-mass term for the neutrino. Such a term, which couples ν_R to the left handed $SU(2)$ doublet neutrino, allows to identify the heavy SM singlet belonging to a N_F-dimensional supermultiplet of the new interaction with a right handed neutrino field. It also lets us recognize its phase as lepton number.

The fact that the theory is derived from an effective four-Fermi interaction is manifested in relations (5.10 - 5.13) since the running couplings approach the corresponding bare couplings as $\mu \to \Lambda$.

A Lagrangian with a canonical kinetic term can be obtained by rescaling the scalar field $\Phi \longrightarrow \Phi/\sqrt{Z_\Phi}$ to get

$$\mathcal{L}_\Phi = \partial_\mu \Phi \partial^\mu \Phi^\dagger + g\left(\bar{\nu}_R^c \nu_R \Phi + \text{h.c.}\right) - V(\Phi). \tag{5.14}$$

In adittion, one can express the theory in terms of physical quantities by means of the redefinition of the bare parameters

$$g = \frac{g_o}{\sqrt{Z}} \tag{5.15}$$

$$m^2 = \frac{m_\Phi^2}{\sqrt{Z}} \tag{5.16}$$

$$\lambda = \frac{\lambda_o}{Z^2}. \tag{5.17}$$

Once this is done, the potential for the scalar field is given by

$$V(\Phi) = m^2 \Phi^\dagger \Phi + \lambda \left(\Phi^\dagger \Phi\right)^2. \tag{5.18}$$

This potential involves only the radial component of the scalar field, i.e. it is symmetric under a global U(1) phase transformation (lepton number). Therefore, if the scalar field acquires a vacuum expectation value

$$v = \sqrt{-\frac{m^2}{\lambda}}, \tag{5.19}$$

breaking spontaneously the U(1) symmetry, the phase field would become a Goldsotone boson, massless at every level in pertubation theory.

However, at energies close to Planck scale, it is expected that any global U(1) symmetry will be broken due to the black-hole dynamics which induces low energy effective operators that do not conserve global charges, such as lepton/baryon number [20]. Thus, we can parametrize the explicit symmetry breaking terms by adding to our Lagrangian the lowest dimension symmetry breaking term that can be constructed out of right handed neutrino fields, [1] i.e.

$$G' \left[(\bar{\nu}_R^c \nu_R)^2 + (\bar{\nu}_R \nu_R^c)^2 \right] . \tag{5.20}$$

This term introduces another unknown high enegy scale Λ', which is inversely proportional to G' , $G' \propto \frac{1}{\Lambda'^2}$, and violates lepton number by four units.

On general grounds a small explicit breaking is expected, such that $\Lambda' > \Lambda$, so that one can also parametrize the effects of the symmetry breaking term by means of the auxiliar scalar field from the compositness condition (5.6)

$$g' \left(\bar{\nu}_R^c \nu_R \Phi^\dagger + \bar{\nu}_R \nu_R^c \Phi \right) . \tag{5.21}$$

With the above expressons in mind, one can derive in a straightforward manner an expression for the mass of the right handed neutrino

$$m_R^2(\theta) = (g^2 + g'^2 + 2gg' \cos(2\theta))v^2, \tag{5.22}$$

where we use the dimensionless parametrization of the angular field $\frac{\phi}{v} = \theta$.

On the other hand, due to the explicitly breaking of the lepton U(1) symmetry, the θ field develops an effective potential from 1-loop corrections which reads

$$V(\theta) = -\frac{1}{(16\pi)^2} m_R^4(\theta) \cdot \ln \left[\frac{m_R^2(\theta)}{v^2} \right] , \tag{5.23}$$

leading to a non-zero mass for the θ field, which becomes now a Pseudo Nambu-Goldstone Boson (PNGB).

[1] This set up introduces, just below the Planck scale, two dimension six operators for ν_R but assumes the absence of the usual dimension three Majorana mass term. This may sound unnatural, however in string theory – our only consistent description of Planck scale physics –, non-generic effective actions below the Planck scale are the natural expectation.

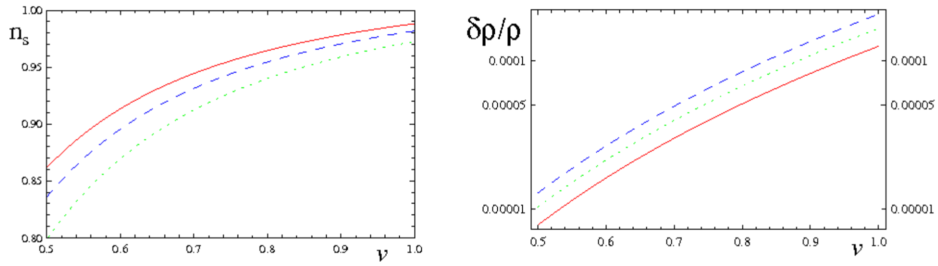

Fig. 5.1: Evolution of the spectral index and the density fluctuations as a function of the spontaneously symmetry breaking scale v for different couplings (g, g').

At this point, it is important to notice two features of this model in relation to its phenomenological behaviour. The first one is that the (true) minimum of the potential is located at $2\theta = \pi$ and does not vanish at it, therefore a redefinition of the potential will be needed. The second one is that since we are assuming a hierarchy between the spontaneous and explicit symmetry breaking scales, being the spontaneous the smallest, i.e $\Lambda' > \Lambda$, the corresponding Yukawa couplings between the scalar field and the neutrinos will exhibit also the same hierarchy, $g' <<< g$.

Taking this into account, the potential for the scalar field, which will drive the inflationary dynamics, takes the following form

$$V(\theta) = M^4 \cdot (1 + \cos(\theta)), \tag{5.24}$$

where M is given in terms of the Yukawa couplings by

$$M^4 = -\frac{g^3 g' v^4}{32\pi^2} \left(1 + 4 \ln g\right). \tag{5.25}$$

The above potential is known in the literature under the name of "Natural Inflation" [21] and for certain range of its parameter space displays a potential flat enough to satisfy the inflation requirements. It also exhibits two widely different energy scales: M which establishes the scale of the potential and will be related with the energy scale at which inflation takes place and v, the vacuum expectation value which will define the mass of the right handed neutrino, that of the inflaton and the scale of spontaneous symmetry breaking, together with the inflationary observables.

In [17], an exhaustive analisis of the inflationary epoch has been performed to constrain the value of the Yukawa couplings and the scale of spontaneous breaking. It was found that

$$0.7\mathrm{M_{Pl}} \leq v \leq 0.9\mathrm{M_{Pl}} \tag{5.26}$$

$$\left(g^3 g'\right)^{1/2} \sim 10^{-5} \tag{5.27}$$

was needed to provide the correct scalar spectral index and size of density fluctuations. Figure (5.1) shows the evolution of these observables as a function of the symmetry breaking scale for different sets of couplings g, g'. As a result, a value of $M^4 \sim (10^{16}\ \mathrm{GeV})^4$ must be enforced for a natural choice of g.

3 Inflationary dynamics

From the potential obtained in last section, one can reconstruct the dynamics of the inflationary field. Tipically, almost any inflationary transition goes trough two recognizible periods. During the first one, the inflaton motion is overdamped by the huge exponential expansion of the Universe, making it evolve very slowly (slow roll phase). Owing to this proccess, the Universe dilutes any undesearible relic and emerges extremely flat and smooth. The second epoch comprises the oscillations of the inflationay field, which gets converted into radiation, "reheating" the Universe. Along this phase, the inflaton mimics nonrelativistic matter evolution.

During the second stage, the decay width of the inflaton can be parametrized as

$$\Gamma \simeq k_0 m_\theta(t), \tag{5.28}$$

where k_0 denotes the coupling between the inflaton and relativistic matter (essentially all particles are massles, i.e. relativistic at that time) and $m_\theta(t) = \sqrt{V''(\theta(t))}$ is the PNGB time variating mass, defined as the second derivative of the potential. The value of k_0, which sets the decay width, determines for how long the inflaton dominates the energy budget of the Universe while reheating.

The equation of motion which governs the dynamics of the inflaton can be read as

$$\ddot{\theta} + (3H + \Gamma)\dot{\theta} + \frac{V'(\theta)}{v^2} = 0, \tag{5.29}$$

where the factor $1/v^2$ arises from the parametrization for the inflaton field we use and Γ is the PNGB decay width operator which takes into account the dilution of the scalar field into radiaton. Contrary to the traditional picture, we include this term even in the inflationary epoch. This is the so-called "warm inflaton" scenario [22]. Strictly speaking one should always include such a term. However, in most of cases, $\Gamma << H$ during the slow roll phase, and one can safely neglect it. In our case, Γ is not so small, so we have to include it at every stage.

The evolution equations for the fields involved are well known and given by

$$\dot{\rho}_\phi = -3H(1 + w_\phi)\rho_\phi - \Gamma v^2 \dot{\theta}^2, \tag{5.30}$$

$$\dot{\rho}_\gamma = -4H\rho_\gamma + \Gamma v^2 \dot{\theta}^2, \tag{5.31}$$

where

$$H^2 = \frac{8\pi}{3M_{Pl}^2}\left(\rho_\theta + \rho_\gamma\right), \tag{5.32}$$

$$\rho_\theta = \frac{1}{2}v^2\dot{\theta}^2 + V(\theta), \tag{5.33}$$

$$V(\theta) = M^4\left(1 + \cos(\theta)\right), \tag{5.34}$$

with the dimensionless parametrization $\theta = \frac{\phi}{v}$.

Solving numerically this set of equations, one obtains the thermal history of the universe and the behaviour of the inflaton during its rolling down of the potential. The general pattern which follows from these equations is clear. The inflaton starts dominating the energy density of the Universe, only diluting away as a consequence of the expansion. Once the end of the slow-roll phase is reached, the friction term becomes dominant and converts the inflaton energy into radiation, reheating the Universe and recovering the old Big Bang picture. The point where both components cross depends on the the precise value

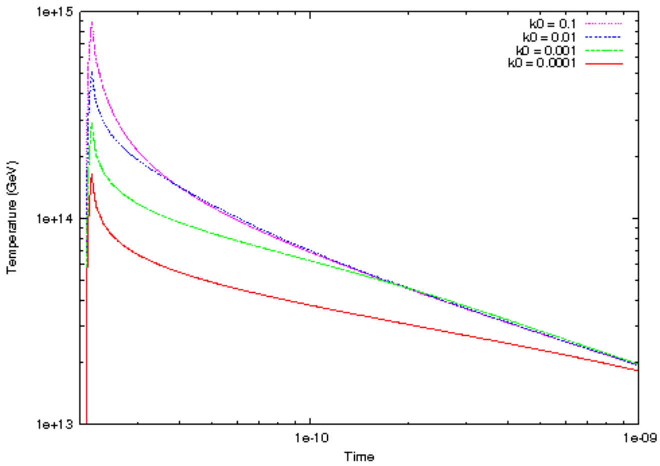

Fig. 5.2: Temperature evolution depending on the reheating parameter k_0.

of Γ. This point signals a time $t \approx \Gamma^{-1}$ at which most of the energy stored in the inflaton field gets converted into radiation. This feature motivates the widely used instantaneous decay approximation, which is a qualitative/easy way to analytically solve the equations of motion and provides a picture that captures the essence of the behaviour of the inflaton during the reheating process.

Alternatively, when one is dealing with an inflationary scheme, it is important to know the temperature reached once the inflaton has completely decayed away. This temperature, normally called the reheating temperature, is determined by the radiation energy density generated as follows

$$T = \left(\frac{30\rho_\gamma}{g_* \pi^2}\right)^{1/4}, \tag{5.35}$$

where g_* is the number of relativistic degrees of freedom in the theory (≈ 100 within the standard model). Fig (5.2) shows the evolution of the temperature for different k_0 values. From there it can be seen that the larger the k_0, the shorter the matter domination period at the end of the inflationary phase.

Between the end of the slow roll phase and the time $t \simeq \Gamma^{-1}$ (the instantaneous decay time), equation (5.31) is dominated by the kinetic term of the θ field and thus, the temperature does not fall as in a radiation-like dominated Universe, but as $T \approx t^{-1/4}$ due to the entropy release of the decays. During this phase, the temperature reaches an almost flat plateaux from the point where the energy density in the radiation born out of the inflaton and the energy density of the inflaton itself became comparable up to $t \simeq \Gamma^{-1}$. This fact can be seen clearly when $k_0 = 0.001$, and $k_0 = 0.0001$ to a greater or lesser extent. After $t \simeq \Gamma^{-1}$, the Universe becomes radiation dominated, the expansion term dominates in eq (5.31) and then, the temperature falls like $T \approx t^{-1/2}$.

Tipically temperatures reached after our inflaton decayed away are of order of 10^{14}GeV. Such high temperatures can be problematic in models beyond the standard model, like supersymmetry, because it would lead to an overproduction of gravitinos [23], which would have catastrophic consequences for the evolution of the Universe and specially in the formation of light elements (H, He..) at BBN [24, 25]. However, they are perfectly acceptable in the context of the standard model.

4 Baryon Asymmetry Calculation

In this section, we calculate the baryon asymmetry generated within this model, for which we derivatively couple the Pseudo Nambu-Goldstone boson to a leptonic current. In order to obtain a non zero expectation value for the time derivative of the Goldstone field, we have included in the Lagrangian a term that soflty breaks the U(1) symmetry explicitly as well as spontaneously. The above mentioned derivative coupling takes the form

$$\frac{1}{f}\partial_\mu \phi J^\mu_L, \tag{5.36}$$

where J^μ_L is the lepton current and f is associated to the energy scale responsible for such a term. This sort of coupling would be only possible if, as happens in our model, lepton asymmetry is violated, otherwise the divergence of the current would vanish. Our inflationary phase now, is just a textbook example of a second order phase transtition, where a scalar order parameter (our phase field) evolves from one field value to another, as the true minimum of its effective potential changes. In the meantime, there will be a period during which the velocity of the field develops an expectation value.

This term implies a Time Reversal and Lorentz invariance violation, which likewise will lead to a temporary violation of CPT. Even though this could scare any responsible reader, to do so locally is perfectly consistent [26]. Mild violations of CPT could have an origin in the neutrino sector [27, 28]. Regarding the possible origin of this term, when dealing with theories near Planck scale, due to non global lepton charge conservation, the divergence of the Lepton current is non-zero, making this term suitable to appear in the Lagrangian as an effective operator.

Regarding baryogenesis, CPT violation in the theory relaxes the Sakharov conditions for generating the baryon asymmetry dynamically. Normally, in addition to the baryon number and CP violation, one has to consider a scenario where thermal equllibrium can not be reached, since along with CPT conservation it enforces the production of a zero net baryon number. The reason is clear (recall that baryon number is an odd quantity under a CPT transformation)

$$\begin{aligned} < \hat{B} > &= \mathrm{Tr}\left[e^{-\beta H_{CPT}} B\right] \\ &= \mathrm{Tr}\left[(CPT)e^{-\beta H_{CPT}}(CPT)^{-1}(CPT)\ \hat{B}\ (CPT)^{-1}\right] \\ &= \mathrm{Tr}\left[(+1)e^{-\beta H_{CPT}} \cdot (-1)\hat{B}\right] = - < \hat{B} > \Longrightarrow < \hat{B} >= 0\,, \end{aligned} \tag{5.37}$$

where H_{CPT} is CPT-conserving Hamiltonian and $\beta = 1/T$. However, the above expression no longer holds when CPT is violated in the Lagrangian and therefore, a net asymmetry can be produced even in thermal equilibrium. In addition, there is no need to break CP (or departure of equilibrium) since the Sahkarov conditions do not apply when CPT is violated .

We are interested in relating the inflationary scalar field with the baryon asymmetry production. Therefore, by identifying $\dot{\phi} \longrightarrow \mu$ and $\frac{v}{f} \longrightarrow \lambda$, a dimensionless coupling to be constrained later, the equation of motion can be re-written as

$$\ddot{\phi} + 3H\dot{\phi} + V'(\phi) + \Gamma\dot{\phi} = -\frac{\lambda}{v}(\dot{n}_L + 3Hn_L), \tag{5.38}$$

where Γ without subscript refers to the usual inflaton decay width into radiation, eq (5.28), and n_L denotes the lepton number density.

The above equations describe the lepton asymmetry produced in thermal equilibrium during the inflaton slow roll down and subsequent decay with the interaction term shown in eq (5.36), provided that the rate of change of $\dot{\phi}$ is sufficiently low. If this were the case, this interaction would shift the lepton and antilepton energy levels like a chemical potential for lepton number. (Here sufficiently low simply means that the typical time scale of lepton violating processes must be fast enough to maintain thermal equilibrium).

If thermal equilibrium cannot be reached, one has to substitute the divergence operator by the operator that violates lepton number. As this term gives rise to the decay of the inflaton field, one can approximate the effect of the decay of the motion of the inflaton field due to its lepton number violating interactions by including an extra friction term, proportional to the width of the lepton number violation

$$\ddot{\phi} + 3H\dot{\phi} + V'(\phi) + \Gamma\dot{\phi} = -\Gamma_L\dot{\phi}\ . \tag{5.39}$$

where Γ_L is the interaction width which violates lepton charge. In this case, we differenciate both, because we are interested in the contribution of this last term to the baryon/lepton asymmetry generated. In general, this term should be included inside the entropy term production Γ.

Comparing both equations, we can see that the Boltzmann equation for the lepton asymmetry is given by

$$\dot{n}_L + 3Hn_L = -\frac{1}{\lambda}\Gamma_L v\dot{\phi}. \tag{5.40}$$

The analitic solution of this Boltzmann equation is given by

$$n_L(t) = n_L(t_o) - \left(\frac{a}{a_o}\right)^3 \times \int_{t_o}^{t_f} dt\frac{1}{\lambda}\Gamma_L v\dot{\phi} \tag{5.41}$$

One notices that the lepton asymmetry is just the area enclosed by the phase field throughout its oscillatory movement around its minimum, during which the inflaton produces leptons/antileptons for positive/negative velocities (by negative we mean velocities in the opposite direction). As this procces is modulated by the expansion of the Universe, it leads to a non zero value.

In our model, the lepton violating operator takes the form

$$\frac{1}{\Lambda'^2}(\bar{\nu}_R^c \nu_R)^2 + h.c, \tag{5.42}$$

This term comes from the explicit lepton number violation term, eq (5.20), with the identification $G' \equiv 1/(\Lambda')^2$, i.e, we are assuming the maximal value this coupling constant may have. Like any dimension six operator, ours yields a decay width of the form

$$\Gamma_L = \lambda^2 \frac{m_\theta^5}{\Lambda'^4} \ , \tag{5.43}$$

where $m_\theta = \sqrt{V''(\theta)}$ is the inflaton mass.

We are interested not in the lepton density or the baryon one, but the baryon to photon ratio. At a temperature T the photon number density is given by

$$n_\gamma = \frac{2\zeta(3)}{\pi^2} T^3, \tag{5.44}$$

Thus, once we solve numerically the equation for the lepton number asymmetry (5.40), we can estimate the baryon asymmetry $\eta = \frac{n_B}{n_\gamma}$. As we have previously mentioned, since this quantity tracks $\dot{\phi}$, one would expect a damped oscillating behaviour asymptotically reaching a final value, once there is no sufficient feedback to keep producing it. We show the particular feature in Fig. (5.3) for the following set of values

$$k_0 = 0.01 \ , \tag{5.45}$$

$$\Lambda' \simeq 10^{16} \text{GeV} \ , \tag{5.46}$$

$$\lambda \simeq 0.01 \ , \tag{5.47}$$

which give the experimentally observed baryon asymmetry. On the other hand, the evolution of the velocity with time will resemble the

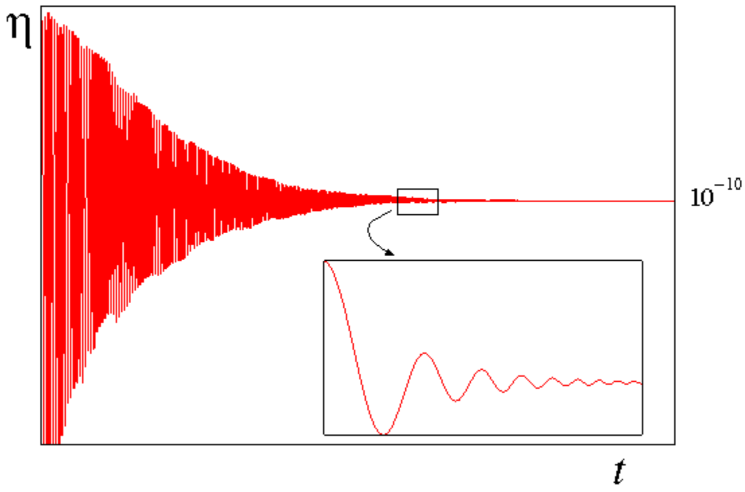

Fig. 5.3: Baryon asymetry evolution for $k_0 = 0.01$, $\lambda = 0,01$ and $\Lambda' = 10^{16}$ GeV in logarithmic scale. The amplified snapshot shows a cartoon picture of the oscillations.

one showed, but asymptotically converging to zero as a consequence of the progresssive lost of kinetic energy.

For this calculation we have taken into account that the lepton to baryon asymmetry convertion by the sphaleron proccess at later epochs is given by

$$n_B = -\frac{28}{79} n_L \ . \tag{5.48}$$

Similarly, as an additional/alternative source of lepton asymmetry we could have used the lepton number violating operator

$$g' \left(\bar{\nu}_R^c \nu_R \Phi^\dagger + \text{h.c} \right) . \tag{5.49}$$

Contrary to the first one, this term is a four dimensional operator, so it would produce a decay width with the following form

$$\Gamma_L = \frac{g'^2}{8\pi} m_\theta . \tag{5.50}$$

Comparing this last term with the one given in eq (5.43), one can make an educated guess for the value of g' from the values of Λ' and λ

needed to get the right amount of baryon asymmetry (eq 5.45 - 5.47). This turns out to be

$$g' \simeq 10^{-7}\,. \tag{5.51}$$

The value of g' can then be used now to determine g by requiring the size of primordial fluctuations to agree with experiment, eq (5.27) and both together constraint the mass of the right handed neutrino m_R (eq 5.22). Consecuently, these parameters take the following values

$$g \simeq 0.1, \tag{5.52}$$

$$m_R \simeq gv = 10^{18}\left(\frac{v}{M_{\rm Pl}}\right)\,\text{GeV}. \tag{5.53}$$

Up to an order of magnitude smaller masses for the right handed neutrino field can be obtained for larger values of g'. However, given that in the $g' < g$ regime we are forced to have $g^3 g' \simeq 10^{-10}$ to provide the right size of scalar density perturbations, $m_R > 10^{16}(v/M_{\rm Pl})$GeV for any choice of fermion couplings.

With this at hand, we can already test our model. As it is well known in addition to scalar (density) perturbations, our field will also give rise to tensor (gravitational wave) perturbations. Generally, the tensor amplitude is given in terms of the tensor/scalar ratio

$$r \equiv \frac{P_T}{P_R} = 16\epsilon \tag{5.54}$$

The tensor to scalar ratio r goes like $g^2 g'^2$ and for our model it turns out to be well below the detection sensitivity of current and (near) future experiment. Gravity waves are the holy grail of next generation of experiments and if found, will rule out this model.

Strictly speaking, n_s is not a constant, and its dependence on the scale can be characterized by its running. Our model predicts a very small and negative spectral index running, scaling as g'/g. It is so negligible small that it is essentially indistinguishable from zero running. Small scale CMB experiments will provide more stringent tests on the running. If these experiments exclude a trivial (consistent with zero) running, i.e. if they detect a strong running, our model would be ruled out.

5 Discussion and Conclusions

In this paper, we have discussed the baryon asymmetry generated in an inflationary model, without a fundamental scalar field. We have showed that it is possible to obtain the observed η value from an inflaton-like composite generated out of strongly coupled right handed neutrinos, while at the same time agreeing with cosmological observations.

The possibility of dynamically generating a scalar field, responsible not only for breaking the symmetry but also for giving mass to the right handed neutrino masses and whose decay generates the baryon asymmetry of the universe by using the CPT non-invariance of the universe during its early history makes the model especially economic and therefore physically appealing.

The resulting model is phenomenologically tightly constrained, and can be experimentally (dis)probed in the near future.

6. ELECTROWEAK BARYOGENESIS WINDOW IN NON STANDARD COSMOLOGIES

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abstract

In this work we show that the new bounds on the Higgs mass are more than difficult to reconcile with the strong constraints on the physical parameters of the Standard Model and the Minimal Supersymmetric Standard Model imposed by the preservation of the baryon asymmetry. This bound can be weakened by assuming a nonstandard cosmology at the time of the electroweak phase transition, reverting back to standard cosmology by BBN time. Two explicit examples are an early period of matter dominated expansion due to a heavy right handed neutrino (see-saw scale), or a nonstandard braneworld expansion.

1 Introduction

A wealth of cosmological observations over the past years have provided a deep knowledge on the thermal history of the Universe since its first nanoseconds, up to today. Supernova candles show that the Universe is now accelerating as a consequence of an exotic particle or more likely a cosmological constant with negative pressure. Measurements from the CMB tell us that our Universe is flat, isotropic and (almost) homogeneous and its physics can be accurately described by the Hot Big Bang Model and General Relativity.

One of the corner-stones of the Hot Big Bang model is Big Bang Nucleosynthesis, the theory about the formation of light elements (namely deuterium, helium, and lithium) that were produced in the first few minutes after the Bang. The abundances of these lights elements depend on the density of protons and neutrons at the time of nucleosynthesis (as these were the only baryons around at this time) and provide a strong evidence for a necessity of a baryon asymmetry, an excess of nucleons over antinucleons. Furthermore, the Universe seems to contain relatively few antibaryons. There is clear evidence that at least the local cluster of galaxies is made of matter, and there is no plausible mechanism to separate matter from antimatter on such large scales.

Then one of the most challenging aspects of the interplay between particle physics and cosmology is to construct a compelling and consistent theory that can explain the observed baryon asymmetry of the universe. The tiny difference between the number density of baryons and antibaryons, of about 10^{-10} if normalized to the entropy density of the Universe. In order to be able to generate such an asymmetry any theory must fulfil certain conditions. These conditions, called the Sakharov's conditions [14] establish the necessary ingredients for the production of a net baryon asymmetry, which are

1. Non conservation of baryon number
2. Violation of C and CP symmetry
3. Departure from thermal equilibrium

The need for the first two conditions is quite obvious. Regarding the third one the Universe must have been out of thermal equilibrium in order to produce net bayon number, since the number of baryons and antibaryons are equal in thermal equilibrium (if B violating processes do exist). It is also important to notice that **all** known interactions are in thermal equilibrium when the temperature of the Universe is between 100 GeV and 10^{12} GeV.

Many mechanisms for the production of the baryon asymmetry have been discussed for different periods of the evolution of the early universe, which include GUT-baryogenesis, leptogensis, etc. Among all the proposals, the generation of the baryon asymmetry at electroweak scale is specially appealing since the electroweak scale is the last instance in the evolution of the Universe in which the baryon asymmetry could have been produced within minimal frameworks. The Standard Model satisfies every Shakharov condition and thus was considered that solely within this framework baryogenesis could be explained.

Firstly, baryon number violation occurs in the Standard Model through anomalous processes. Secondly, at low temperatures this anomalous baryon number violation only proceeds via tunnelling which is exponentially suppressed. However, anomalous baryon number violation is rapid at high temperatures and the weak phase transition, if first order with supercooling, provides a natural way for the Universe to depart from equilibrium at weak scale temperatures. Electroweak phase transition can be then seen as bubbles of the broken phase which expand and end up filling the Universe. In this picture, local departure takes place in the vicinity of these expanding bubble walls. Lastly, C and CP are known to be violated by the electroweak interactions. So, in principle, all the required ingredients are there.

However, the standard model fails in almost every aspect. The CKM phase, the only source for CP violation in the standard model, is extremely small to explain the observed baryon to entropy ratio. Another decisive check comes from the requirement that any net baryon asymmetry produced during the transition should survive until today. For an Universe whose expansion rate is slower than the anomalous baryon violating processes, thermal equilibrium would be recovered after the electroweak phase transition. Therefore, any asymmetry in baryon number created during the transition would be erased. In the broken phase, the rate of baryon number violation is exponentially suppresed by a factor $\mathcal{O}\left(\phi/gT\right)$, where ϕ is the value of the order parameter and g is the weak coupling constant. Thus, when demanding the baryon violating width to be smaller than Hubble rate, one finds

$$\frac{\phi(T_{ew})}{T_{ew}} \gtrsim 1\,, \tag{6.1}$$

where T_{ew} stands for the temperature at which the electroweak phase transition is completed. Usually this temperature can be safely approximated to the critical temperature T_c when both phases co-exist. The above condition constitutes the so called "sphaleron bound", and can give new information and constraints about the CP and Higgs sectors of the Standard Model. In particular, it has been shown that Higgs masses larger than 40 GeV can be ruled out by imposing that the baryon asymmetry of the Universe be generated during the weak transition [29, 30].

Nonetheless, the sphaleron bound (eq. 6.1) presented above, assumes a particular thermal history of the Universe, one where during the electroweak phase transition the energy density of the universe was dominated by radiation. In section 2, we will show that, under different thermal histories of the Universe or different cosmologies, a less stringent condition can be obtained, permitting Higgs masses above the current experimental bounds. In section 3, we will analyse a scenario with a non standard thermal history during the electroweak phase transition which leads to a modified sphaleron bound condition, while in section 4 we relax this bound by modifying the underlying cosmology. We will conclude in section 5.

2 Sphaleron Bound reviewed

The evolution of any baryon asymmetry in comoving units during the electroweak phase transition can be written as

$$\frac{n_{\rm freeze}}{n(t_B)} = \exp\left[-\int_{t_b}^{\infty} dt\;\widetilde{\Gamma}_{{\rm sph}(t)}\right]\,, \tag{6.2}$$

where n_{freeze} is the baryon asymmetry which survives to partake of nuclesynthesis, $n(t_B)$ is the baryon asymmetry at the beginning of the phase transition and t_b is the time at which the bubble nucleation proceeds, starting up the phase transition.

The meaning of this equation is clear. The baryons created at the bubble walls are subject to decay after they enter the broken phase, if the baryon number violating processes are not sufficiently suppressed. We should require then this attenuation not to reduce the created asymmetry to less than that required for nuclesynthesis *i.e.*

$$\int_{t_b}^{\infty} dt\ \widetilde{\Gamma}_{\text{sph}(t)} = -\log\left(\frac{n_{\text{freeze}}}{n(t_B)}\right) \leq 1\ . \tag{6.3}$$

The sphaleron width is given by [31]

$$\widetilde{\Gamma}_{\text{sph}}(t) = \alpha_n 6 N_F^2\ \mathcal{C}\ g\frac{\phi^7}{T^6}e^{-\frac{E_{\text{sph}}}{T}}\ , \tag{6.4}$$

where α_n is a number of order one, whose precise value depends on the model and its corresponding set of conserved charges and N_F is the number of fermion families. $\mathcal{C}$ is a temperature independent parameter accounting for the degrees of freedom of the sphaleron and may be expressed in the following way

$$\mathcal{C} = \left(\frac{\omega_-}{2\pi g\phi(T)}\mathcal{N}_{\text{tr}}\mathcal{N}_{\text{rot}}\mathcal{V}_{\text{rot}}\mathcal{K}_{\text{sph}}\right)\ . \tag{6.5}$$

where ω_- is the frequency of the negative mode of the sphaleron, $\mathcal{V}_{\text{rot}} = 8\pi^2$, $\mathcal{N}_{\text{tr}}\mathcal{N}_{\text{rot}} \simeq 86 - 5\ln(m_H^2/8m_W^2)$ and $\mathcal{K}_{\text{sph}} = \{7.54, 5.64, 4.57, 3.89, 3.74\}$ for $m_H = \{0.4, 0.5, 0.6, 0.8, 1\}m_W$, and extrapolated for other values of m_H.

As the dominant contribution to the integral (6.2) comes from temperatures very close to T_B, it can be approximated to its value at this temperature. Such approximation slightly overestimates the dilution.

As $t \sim H^{-1}$, this yields the condition

$$\widetilde{\Gamma}_{sph}(t_b) \leq H(t_b)\ . \tag{6.6}$$

This equation shows what we pointed out before, the sphaleron rate processes must be slow enough, i.e out of thermal equilibrium, in order that any $(B+L)$ asymmetry won't be erased. This bound is usually stated as a lower bound on the sphaleron energy, or as a lower bound on the ratio of the vev to the temperature at the critical temperature and can then be converted into a bound on the parameters in a specific model.

Usually the literature shows this bound in the conventional cosmological scenario, that is, in a radiation dominated Universe. Within this scenario the expansion rate is given by

$$H^2_{\rm rad} = \frac{4\pi^3}{45M^2_{Pl}} g_* T^4 \ . \tag{6.7}$$

Inserting this expression in eq. 6.6, one finds that

$$\frac{\phi_c}{T_c} \gtrsim \frac{1}{\mathcal{B}}\sqrt{\frac{4\pi}{\alpha_w}} \left(7\log\frac{\phi_c}{T_c} + \log W(T_c) - \log H_{\rm rad}\right) \ , \tag{6.8}$$

where $\mathcal{B} = \{1.52, 1.61, 1..83, 2.10\}$ for $m^2_H/m^2_W \in \{0.008, 0.08, 0.8, 8\}$ and quadratically interpolated for intermediate values and $W(T) = 6\alpha_n N^2_f \mathcal{C} g T_c$. Solving this equation numerically gives

$$\frac{\phi_c}{T_c} \gtrsim 1 \ . \tag{6.9}$$

Alternatively this bound can be restated as a function/bound on different cosmological scenarios for which the expansion rate takes a different value. In such scenarios [32]

$$\frac{\phi_c}{T_c} \gtrsim \frac{1}{\mathcal{B}}\sqrt{\frac{4\pi}{\alpha_w}} \left(7\log\frac{\phi_c}{T_c} + \log W(T_c) - \log H_{\rm rad}\right) + \delta_{\frac{\phi_c}{T_c}} \ , \tag{6.10}$$

where

$$\delta_{\frac{\phi_c}{T_c}} = \frac{1}{\mathcal{B}}\sqrt{\frac{4\pi}{\alpha_w}} \log\frac{H}{H_{\rm rad}} \ . \tag{6.11}$$

This new term has the effect of relaxing the sphaleron bound. This effect can be seen in figure 6.1, where the difference between the solutions given by eq. 6.8 and 6.10, *i.e,* $\Delta\left(\frac{\phi_c}{T_c}\right) = \frac{\phi_c}{T_c}\Big|_{H_{rad}} - \frac{\phi_c}{T_c}\Big|_{H}$ is plotted for different values of H.

In addition, it is clear that only drastic modifications, *i.e.* modifications where the energy density (and therefore the expansion rate) is several orders of magnitude larger than the one given in a radiation dominated scenario, can relax the bound in a sensible way. We are interested in studying whether such a modification to the sphaleron bound can be helpful to open up the allowed parameter space for electroweak baryogenesis. To study this, let us review first how this bound is obtained in the Standard Model, and what its implications are.

In the Standard electroweak theory the effective potential at high temperatures reads as [33]

$$V(\phi, T) \approx \frac{M(T)^2}{2} - ET\phi^3 + \frac{\lambda_T}{4}\phi^4 \,, \tag{6.12}$$

where $M(T), B$ and λ_T are the temperature dependent effective mass, cubic term and quartic coupling respectively; given at the one-loop ring improved values

$$\begin{aligned}
M(T) &= \sqrt{A(T^2 - T_0^2)}\,, \\
A &= \frac{2m_W^2 + m_Z^2 + 2m_t^2}{4v^2} + \frac{1}{2}\lambda_T\,, \\
E &= \frac{2}{3}\left(\frac{1}{2\pi}\frac{2m_W^3 + m_Z^3}{v^3} + \frac{1}{4\pi}\left(3 + 3^{\frac{3}{2}}\right)\lambda_T^{\frac{3}{2}}\right), \\
\lambda_T &= \frac{m_H^2}{2v^2} - \frac{3}{16\pi^2 v^4}\left(2m_W^4 \ln\frac{m_W^2}{a_B T^2} + m_Z^4 \ln\frac{m_Z^2}{a_B T^2} - 4m_t^4 \ln\frac{m_t^2}{a_F T^2}\right), \\
T_0^2 &= \frac{m_H^2 + 8\beta v^2}{2A}\,, \qquad \beta = \frac{3}{64\pi^2 v^4}\left(4m_t^4 - 2m_W^4 - m_Z^4\right).
\end{aligned} \tag{6.13}$$

where T_0 is the temperature at which the phase transition ends, v = 246GeV is the usual Higgs vacuum expectation value at zero temperature, $a_B = (4\pi)^2 e^{-2\gamma_E} \simeq 50$, $a_F = (\pi)^2 e^{-2\gamma_E} \simeq 3.1$, and γ_E is Euler's constant.

By minimizing the effective potential one finds that the ratio of the temperature dependent Higgs vacuum expectation value to the

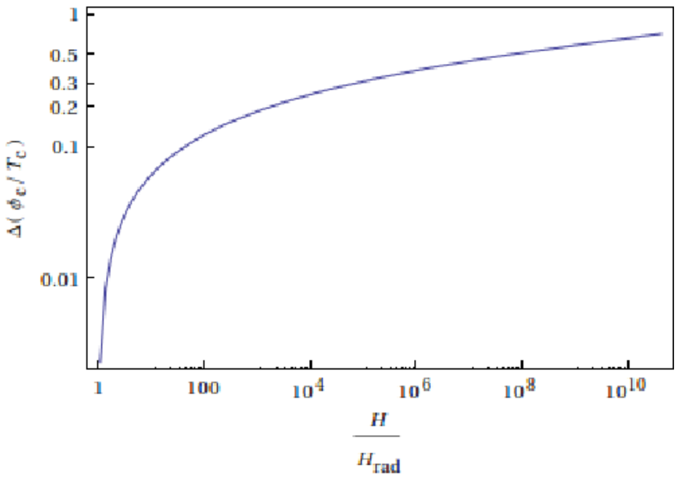

Fig. 6.1: Dependence of the relaxation of the sphaleron bound on the Hubble rate

temperature at a temperature at which a new degenerate minimum appears (the critical temperature) is given by

$$\frac{\phi}{T} = \frac{B + \sqrt{B^2 - 4\lambda_T A(1 - \frac{T_0^2}{T^2})}}{2\lambda_T} \, . \tag{6.14}$$

Using this result as a constraint on the model we can conclude that in order to have a sufficiently strong phase transition within the Standard Model the higgs mass should be smaller than 40 GeV, in clear contradiction with current observations. This is why the Standard Model fails to accommodate a mechanism to generate the baryon asymmetry during the electroweak phase transition. Nevertheless, we have already showed that the sphaleron bound could be weakened by resorting to alternative thermal histories with significantly different expansion rates at the electroweak phase transition. We have yet to see whether the relaxation obtained can be large enough to allow current bounds on the Higgs masses.

There are in the literature plenty of extensions to the standard scenario with an enlarged matter sector where new effects appear and give rise to an enhancement of the strength of the phase transition.

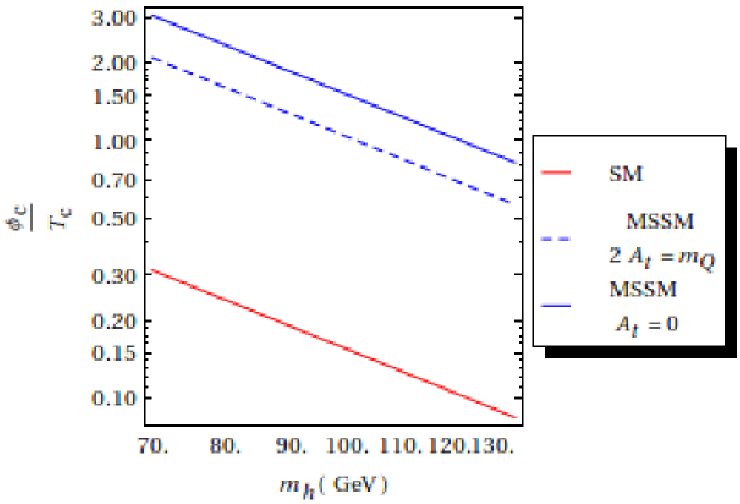

Fig. 6.2: Dependence of the ratio ϕ_c/T_c which controls the preservation of the baryon assymetry on the Higgs mass (treated here as a free parameter) for several models

One of the most interesting extensions is supersymmetry and the so-called "light stop scenario" [34, 35]. In such scenario, stops are light enough, compared to the rest of superpartners, to affect the trilinear coupling to the higgs potential (finite corrections from heavy particles are highly suppressed). This effect impacts the ratio of the temperature dependent vev to the critical temperature in the following way

$$\left.\frac{\phi_c}{T_c}\right|_{MSSM} = \left.\frac{\phi_c}{T_c}\right|_{SM} + \frac{2m_t^3}{\pi v m_h^2}\left(1 - \frac{\tilde{A}_t^2}{m_Q^2}\right)^{\frac{3}{2}} , \tag{6.15}$$

with

$$\tilde{A}_t = A_t - \mu/\tan\beta$$

the effective stop mixing parameter and m_Q the soft supersymmetry breaking mass term for the stops. We can easily see that zero mixing makes the phase transition stronger so a parameter space for this mixing close to zero is highly favoured. However, the mixing to the stops has also an important effect on the one loop corrections to the Higgs mass

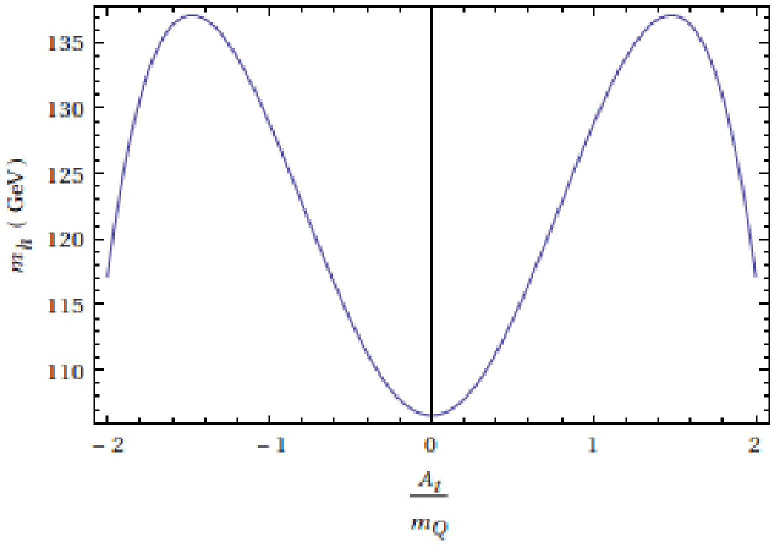

Fig. 6.3: Dependence of the Higgs Mass on the stop mixing. The stop mass $m_{\tilde{t}_l}$ is fixed and given approximately by the soft supersymmetry breaking mass term m_Q

(in the decoupling limit, $M_A >> M_Z$, and a strong hierarchy in the stops spectrum) [36]

$$m_h^2 = M_Z^2|\cos 2\beta|^2 + \frac{3m_t^4}{4\pi^2 v^2}\left\{\log\frac{m_{\tilde{t}_r}^2 m_{\tilde{t}_l}^2}{m_t^4} + \frac{A_t^2}{m_{\tilde{t}_l}^2}\left[2\left(1+\frac{m_{\tilde{t}_r^2}}{m_{\tilde{t}_l^2}}\right) - \frac{A_t^2}{m_{\tilde{t}_l}^2}\left(1+4\frac{m_{\tilde{t}_r^2}}{m_{\tilde{t}_l^2}}\right)\right]\log\frac{m_{\tilde{t}_l^2}}{m_{\tilde{t}_r^2}} + 2\frac{A_t^4}{m_{\tilde{t}_l}^4}\left(1+2\frac{m_{\tilde{t}_r^2}}{m_{\tilde{t}_l^2}}\right)\right\} \tag{6.16}$$

Therefore, while non zero /strong mixing enhances the Higgs mass, it does have the opposite effect on the strength of the electroweak phase transition. We can see this behaviour in the figures 6.2 and 6.3. So even in extensions to the standard scenario, a relaxation on the sphaleron bound would be welcome.

3 First order phase transition in a matter dominated universe

We have seen in the previous section that it is possible to make the electroweak phase transition strong enough to avoid the erasement of the asymmetry by sphalerons if the value of the expansion rate at that scale is orders of magnitude larger than the expansion in the standard, radiation dominated, scenario. In the following, we will study a scenario where this condition can be naturally achieved.

Because the Universe was extremely hot during its early stages, all kind of interesting particles (some yet to be discovered, some which hasn't even been postulated) were present in significant amounts. For $T \gg m$, the mass of the particles in question, their equilibrium abundance is, to within numerical factors, equal to that of photons. When the temperature of the thermal bath drops below m, the equilibrium abundance of such particles is less than that of photons and their contribution to the total energy density becomes suppressed by a factor,

$$(m/T)^{5/2} \exp^{-m/T} , \tag{6.17}$$

except if one (or more) of such particles, which in the following we will call X, drops out of equilibrium and its abundance freezes out (we are assuming that X annihilation cross section is very suppressed). In this case, the relic abundance of X relative to photons remains approximately constant and the contribution to the energy density of X grows as $1/T$ as compared to that of photons. It is obvious then, that eventually the energy density of X will dominate that of the Universe. If the X particle is unstable (but long lived enough) and decays into relativistic particles which thermalise (releasing large amounts of entropy) the Universe will re-enter a radiation dominated era. This will be the scenario we will focus on.

If we assume a flat Universe (as given by observations) the evolution

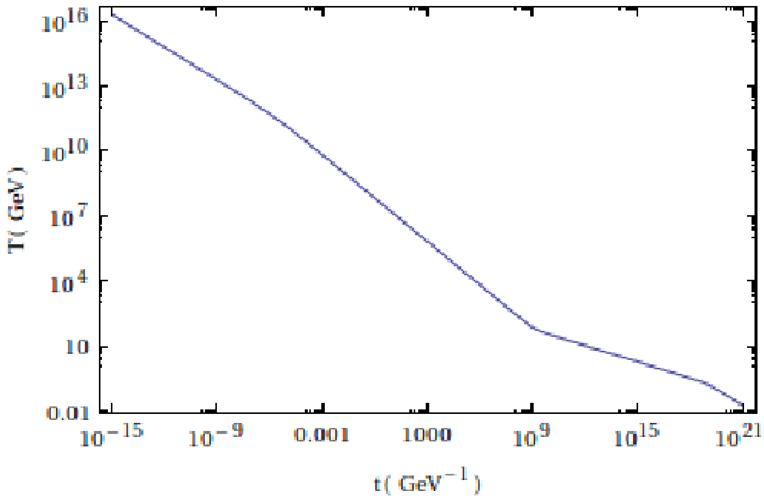

Fig. 6.4: Evolution of the Temperature in an Universe which goes through different epochs

equations for the different components of the Universe are given by

$$\dot{\rho}_X = -3H\rho_X - \Gamma_X \rho_X \tag{6.18}$$

$$\dot{\rho}_r^{\text{old}} = -4H\rho_r^{\text{old}} \tag{6.19}$$

$$\dot{\rho}_r^{\text{new}} = -4H\rho_r^{\text{new}} + \Gamma_X \rho_X \tag{6.20}$$

$$H^2 = \frac{8\pi}{3M_{Pl}^2}\left(\rho_X + \rho_r^{\text{old}} + \rho_r^{\text{new}}\right) , \tag{6.21}$$

where ρ_X is the energy density associated to the particle X, once it becomes nonrelativistic and Γ_X its decay width, ρ_r^{old} is the energy density in radiation not associated with X decays, while ρ_r^{new} is the one coming from X decays [1].

Contrary to the standard picture, the temperature of this universe,

[1] Evolution equations with tracking, *i.e.* when the different components of the Universe chase each others abundance, can also produce early periods of matter domination [37]

will have two sources

$$T(t) = \left(\frac{30}{\pi^2 g_*} \left(\rho_r^{\text{old}}(t) + \rho_r^{\text{new}}(t) \right) \right)^{1/4} \tag{6.22}$$

and therefore its temperature profile, shown in figure 6.4, will be significatly different from that of the standard case [38]. We will start at temperatures larger than the mass of our particle X, M_X, with a radiation dominated universe, where X is in thermal equilibrium. During this time the temperature scales like $t^{-1/2}$ (or $1/a$, being a the scale factor). Once the temperature drops to

$$T_{\text{start}} = \frac{4}{3} r M_X \,, \tag{6.23}$$

with $r = g_X/2$ if X is a boson and $r = 3g_X/8$ if it is a fermion, being g_X the total number of spin degrees of freedom of X, we enter a matter dominated period. During the first part of this period, which comprises most of the matter dominated era, and although X is decaying through an exponential law

$$\rho_X \simeq \frac{2\pi^2 g_*}{45} r M_X T^3 e^{-\Gamma_X t} \,, \tag{6.24}$$

the exponential factor does not affect in a significant way X abundance, the radiation released by X decays is negligible compared with that not coming from X decays, and the temperature falls as in a pure matter dominated period, *i.e.* $T \propto t^{-2/3} \propto 1/a$.

As t approaches $1/\Gamma_X$, the new radiation starts to be comparable with the old one. Thus, X quickly dissapears into (new) radiation and $T \propto t^{-1/4} \propto 1/a^{3/8}$. Once the age of Universe exceeds $1/\Gamma_X$, our matter dominated Universe turns into a radiation dominated one and the temperature starts once more to track the scale factor $T \propto t^{-1/2} \propto 1/a$. At this point

$$T_{\text{end}} = 0.78 g_*^{-1/4} \sqrt{M_{Pl} \Gamma_X} \,. \tag{6.25}$$

The described stages the Universe goes through are depicted in figure 6.5, for a particular choice of parameters. Since we want to recover

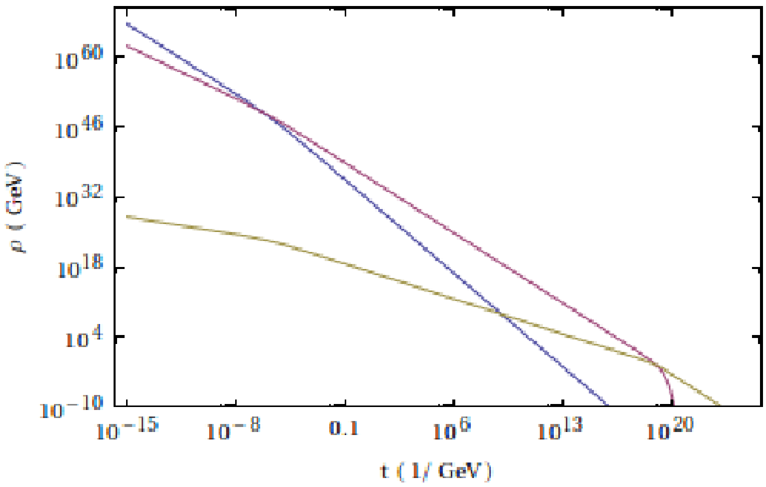

Fig. 6.5: Evolution for the different components of the Universe. Purple, blue and brown line correspond to $\rho_X, \rho_r^{\text{old}}, \rho_r^{\text{new}}$ respectively. In this case, the decay width is $\Gamma_X = 10^{-19}$ GeV^{-1}

the standard (radiation dominated) picture before nucleosynthesis, we can derive a lower bound on Γ_X by requesting that the Universe must have left at the latest the matter dominated era shortly before BBN, which reads

$$\Gamma_{\text{X}} \geq 2.0 \cdot 10^{-24} \sqrt{\frac{g_*}{200}} \text{ GeV} . \tag{6.26}$$

As stated in the previous section an essential condition for electroweak baryogenesis is that the sphaleron transitions would be turned off after the phase transition so that no washing out of the asymmetry produced during the transition occurs. This situation is achieved when the transition rate of the sphaleron interactions is small as compared to the Hubble rate *i.e.* , when these transitions are out of equilibrium. In a radiation dominated universe, the Hubble parameter scales like $H \propto T^2$, however in a matter dominated one, during its first period, when the decays does not significantly reduce the abundance of the X particle

$$H_{\text{MD}}^2 = \frac{16\pi^3 g_* r}{135} \left(\frac{M_{\text{X}} T^3}{M_{Pl}^2} \right) . \tag{6.27}$$

It is then straightforward to notice that large $M_{\rm X}$ masses can substantially change the value of the Hubble rate at electroweak scale and consequently affect the strength of the phase transition.

In order to quantify this variation, we must distinguish between two cases:

(i) The electroweak phase transition temperature is reached when the temperature is essentially given by radiation not coming from X decays, *i.e.*

$$\rho_r^{\rm old}\Big|_{\rm EW} \gg \rho_r^{\rm new}\Big|_{\rm EW}, \qquad (6.28)$$

where in terms of the model parameters

$$\rho_r^{\rm new} \simeq 0.221\, \Gamma_X\, M_{Pl}\, \sqrt{g_* T_{\rm EW}^3 M_X} \qquad (6.29)$$

implies that in this scenario we require

$$\Gamma_X \ll \frac{2.46\times 10^{-13}}{\sqrt{r}}\sqrt{\frac{({\rm GeV})}{M_{\rm X}}}\,. \qquad (6.30)$$

In this case, the extra contribution to the sphaleron bound is given by

$$\delta_{\frac{\phi_c}{T_c}} = \frac{1}{2\mathcal{B}}\sqrt{\frac{4\pi}{\alpha_w}}\log\frac{4r}{3}\frac{M_X}{T_{\rm EW}} \qquad (6.31)$$

so that it gets modified as

$$\Delta\left(\frac{\phi_c}{T_c}\right) \approx -\frac{\log\frac{4r}{3}\frac{M_X}{T_c}}{\frac{1}{\mathcal{B}}\sqrt{\frac{4\pi}{\alpha_w}} - 7\frac{1}{\left(\frac{\phi_c}{T_c}\right)}}\,. \qquad (6.32)$$

(ii) The temperature at the phase transition is set by the radiation coming from X decays, *i.e.*

$$\rho_r^{\rm new}\Big|_{\rm EW} \gg \rho_r^{\rm old}\Big|_{\rm EW} \qquad (6.33)$$

so that

$$\frac{\rho_X}{\rho_r^{\rm new}} \simeq \frac{g_*}{1.9\times 10^{33}}\left(\frac{T_{\rm EW}}{\Gamma_X}\right)^2 \qquad (6.34)$$

and the relaxation on the sphaleron bound reads

$$\Delta\left(\frac{\phi_c}{T_c}\right) \approx -\frac{\log \frac{g_*}{1.9\times 10^{33}}\left(\frac{T_{\rm EW}}{\Gamma_X}\right)^2}{\frac{1}{\mathcal{B}}\sqrt{\frac{4\pi}{\alpha_w}} - 7\frac{1}{\left(\frac{\phi_c}{T_c}\right)}} \; . \tag{6.35}$$

From these equations it is clear that although both schemes can significantly relax the sphaleron bound, they give rise to different phenomenological scenarios. We will come back to this point again later.

But this is not the end of the story regarding the consequences of an early period of matter domination. As it is well known, an early period of matter domination, triggered by a super heavy unstable but longlived particle which goes out of equilibrium at early times and comes to dominate the energy density of the Universe, leads to a reduction of the required number of e-folds before the end of inflation at which the scales of interest today left the horizon. This reduction, which relaxes the flatness condition for the inflationary potential, is due to the fact that the comoving horizon scale grows as $a^{1/2}$ during a matter dominated epoch in contrast to the radiation dominated one where the comoving horizon grows as a. As a consequence, the longer the period of matter domination, the smaller the growth of the universe from the end of inflation up today and therefore the smaller the number of efolds required. This reduction is given by [8]

$$\Delta N = \frac{1}{4}\log\left(\frac{a_{\rm end}}{a_{\rm start}}\right) \; , \tag{6.36}$$

where $a_{\rm end}$ and $a_{\rm start}$ are the scale factor at the beginning and end of the matter dominated era respectively. In terms of the parameters which define our model, *i.e.* the mass and decay width of X, this ratio between the scale factors reads

$$\frac{a_{\rm end}}{a_{\rm start}} = \left(\frac{H_{\rm start}}{H_{\rm end}}\right)^{2/3} \approx 3.9\left(\frac{g_*}{\Gamma_{\rm X} M_{Pl}}(rM_{\rm X})^4\right)^{1/3} \; . \tag{6.37}$$

For values of the decay width close to its lower bound and masses of the order of 10^{15} GeV, this reduction turns out to be over 10 e-foldings.

Likewise, the end of a matter dominated Universe driven by the decay of a long lived massive particle leads to an important entropy production

$$\frac{S_{\text{end}}}{S_{\text{start}}} = \left(\frac{T_{\text{end}} a_{\text{end}}}{T_{\text{start}} a_{\text{start}}}\right)^3 = \left(\frac{\frac{a_{\text{end}}}{a_{\text{start}}}\Big|_{\text{MD}}}{\frac{a_{\text{end}}}{a_{\text{start}}}\Big|_{\text{RD}}}\right)^3 \simeq 12.2 \left(\frac{g_*}{\Gamma_{\text{X}}^2 M_{Pl}^2}\right)^{\frac{1}{4}} (r M_{\text{X}}) \ . \tag{6.38}$$

As it is well known, supersymmetry as well as most of the theories beyond the Standard Model are riddled with new particles associated to new (and higher) energy scales which produce undesirable relics whose abundances, or they mere presence at certain times, do not agree with the current experimental observations of our universe, *e.g.* moduli and gravitinos. So a large release of entropy might help to dilute them, softening (or completely erasing) the constraints on their masses. Consequently, this scenario provides the same services as thermal inflation (regarding the unwanted relics) but without introducing another scalar particle into the theory [9].

On the other hand, it is also important to note that the entropy production that can so nicely solve the unwanted relic problem, can also erase the baryon asymmetry produced at the electroweak scale. Such erasement is given by

$$\eta = \eta_{\text{EW}} \left(\frac{S_{\text{end}}}{S_{EW}}\right) \ , \tag{6.39}$$

where $\eta_{\text{EW}} \approx n_B/s$ is the baryon to photon ratio produced at the electroweak scale and the entropy is given by $S = g_* a^3 T^3$.

As mentioned before, at late times into the matter dominated period $a \propto 1/T^{8/3}$ and then

$$\eta = \eta_{\text{EW}} \left(\frac{T_{\text{EW}}}{T_{end}}\right)^5 \ , \tag{6.40}$$

which in terms of the model reads

$$\eta = \eta_{\text{EW}} \frac{1.5 \times 10^{42}}{g_*^{5/4}} \left(\frac{\Gamma_X}{T_{EW}}\right)^{5/2} \ . \tag{6.41}$$

It is thus clear that we need to generate a large baryon to photon ratio, a ratio of order one or even larger. This needs that the mechanism for baryogenesis to be orders of magnitude more efficient that the standard case, something clearly difficult but not impossible.

One can also see that the entropy production is directly proportional to the decay width, the larger the decay width, the less restrictive the erasement becomes, one would be tempted then to push into the large Γ_X regime. However, large decay widths lead us to scenario (ii) where the temperature is essentially given by the radiation coming from X decays, a scenario where the relaxation of the sphaleron bound is inversely proportional to the decay width. So any gain in the relaxation of the sphaleron bound means a loss in the asymmetry produced. This tension between both scenarios may be seen explicitly in the tables of figure 6.6. Consequently it is clear that the “optimal" case, where we maximize the relaxation of the sphaleron bound and at the same time minimize the dilution of the asymmetry occurs when

$$\rho_r^{\rm old}\Big|_{\rm EW} \approx \rho_r^{\rm new}\Big|_{\rm EW} . \tag{6.42}$$

In this case, the sphaleron bound, for a broad range of M_X values, weakens to

$$\frac{\phi_c}{T_c} \gtrsim [0.64 - 0.69] , \tag{6.43}$$

which may be sufficient to open the window to electroweak baryogenesis in many extensions of the SM and particularly in the MSSM. Figure 6.7 shows the ratio of the temperature dependent vev at the critical temperature to the critical temperature for different Higgs masses as a function of the stop mass (each pair Higgs-stop mass determines the corresponding mixing). The shadowed region signals the reduction that can be obtained for a range of masses and decay widths characterizing the longlived but unstable particle X from the usual $\phi_c/T_c > 1$ bound for preservation of the asymmetry in the standard cosmological scenario. From there it can be clearly seen that a Higgs on the 125-135 GeV range could be made compatible with electroweak baryogenesis, if the thermal history of our universe includes a prolongued period of matter domination.

		$M_X = 10^{12}$ GeV				
		$\rho_{\rm X}^{\rm ew}$ (GeV)	$\rho_{\rm r,old}^{\rm ew}$ (GeV)	$\rho_{\rm r,new}^{\rm ew}$ (GeV)	T_d (GeV)	$\Delta\,(\phi_c/T_c)$
Γ_X (GeV)	10^{-15}	$3.47 \cdot 10^{13}$	1.61	$9.94 \cdot 10^{9}$	17.10	0.10
	10^{-17}	$3.48 \cdot 10^{17}$	$3.48 \cdot 10^{5}$	$9.94 \cdot 10^{9}$	1.71	0.23
	10^{-19}	$5.29 \cdot 10^{20}$	$6.08 \cdot 10^{9}$	$3.88 \cdot 10^{9}$	0.17	0.33
	10^{-20}	$7.37 \cdot 10^{20}$	$9.46 \cdot 10^{9}$	$4.58 \cdot 10^{8}$	0.05	0.34
	10^{-21}	$7.61 \cdot 10^{20}$	$9.87 \cdot 10^{9}$	$4.65 \cdot 10^{7}$	0.01	0.34
	10^{-22}	$7.64 \cdot 10^{20}$	$9.93 \cdot 10^{9}$	$4.66 \cdot 10^{6}$	$5.41 \cdot 10^{-3}$	0.34
	10^{-23}	$7.64 \cdot 10^{20}$	$9.93 \cdot 10^{9}$	$4.66 \cdot 10^{5}$	$1.71 \cdot 10^{-3}$	0.34
	10^{-24}	$7.65 \cdot 10^{20}$	$9.93 \cdot 10^{9}$	$4.66 \cdot 10^{4}$	$5.41 \cdot 10^{-4}$	0.34

		$M_X = 10^{15}$ GeV				
		$\rho_{\rm X}^{\rm ew}$ (GeV)	$\rho_{\rm r,old}^{\rm ew}$ (GeV)	$\rho_{\rm r,new}^{\rm ew}$ (GeV)	T_d (GeV)	$\Delta\,(\phi_c/T_c)$
Γ_X (GeV)	10^{-15}	$3.48 \cdot 10^{13}$	$1.61 \cdot 10^{-4}$	$9.94 \cdot 10^{9}$	17.11	0.11
	10^{-17}	$3.48 \cdot 10^{17}$	34.76	$9.94 \cdot 10^{9}$	1.71	0.23
	10^{-19}	$3.48 \cdot 10^{21}$	$7.49 \cdot 10^{6}$	$9.94 \cdot 10^{9}$	0.17	0.36
	10^{-20}	$2.25 \cdot 10^{23}$	$1.94 \cdot 10^{9}$	$7.99 \cdot 10^{9}$	0.05	0.42
	10^{-21}	$6.81 \cdot 10^{23}$	$8.52 \cdot 10^{9}$	$1.39 \cdot 10^{9}$	0.02	0.44
	10^{-22}	$7.58 \cdot 10^{23}$	$9.82 \cdot 10^{9}$	$1.47 \cdot 10^{8}$	$5.41 \cdot 10^{-3}$	0.44
	10^{-23}	$7.65 \cdot 10^{23}$	$9.93 \cdot 10^{9}$	$1.47 \cdot 10^{7}$	$1.71 \cdot 10^{-3}$	0.44
	10^{-24}	$7.65 \cdot 10^{23}$	$9.93 \cdot 10^{9}$	$1.47 \cdot 10^{6}$	$5.41 \cdot 10^{-4}$	0.44

Fig. 6.6: Value of the important parameters taking special role in our particular scenario dominated by a heavy particle with mass $M_X = 10^{12}$ and 10^{15} GeV and for a set of different decay widths Γ_X. T_d stands for the recovery point of the common radiation dominated era. We have assumed that the thermal history of the Universe begun at about $T \sim 10^{17}$ GeV.

At this point, we must discuss if a particle exists with the characteristics described above. We are looking for a super heavy particle, with an extremely long lifetime in thermal equilibrium at temperatures above its mass. The only particle that appears in (almost) all the extensions of the Standard Model that fulfils these requirements is beyond any doubt the right handed neutrino. Right handed neutrinos through the see-saw mechanism are the fine-tuning-free minimal extension of the Standard Model able to reproduce the only evidence we have observed so far beyond the Standard Model, the light neutrino masses (if their Yukawa couplings are small enough).

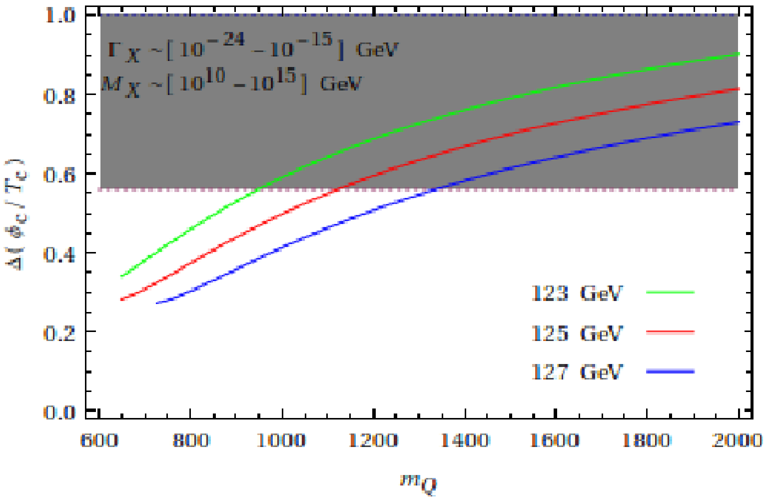

Fig. 6.7: Dependence on ratio of the temperature dependent vev to the temperature at T_c on the stop mass ($\sim m_Q$) for a fixed Higgs mass. The shadowed region shows the relaxation on the sphaleron bound that can be obtained by an erly period of matter domination triggered by a long lived but unstable heavy particle

Of course there are not one but three right handed neutrinos, and their mass matrices and Yukawa couplings are strictly model dependent. However, in a fairly model independent way the mass and lifetime of our X particle, if a right handed neutrino, satisfies

$$\Gamma \propto \frac{m_i M_i^2}{v^2} \tag{6.44}$$

being m the observed light neutrino mass of flavour i, which implies a hierarchycal scenario with an negligible small lightest mass (indistinguishable from zero from an experimental point of view).

4 First order phase transition in braneword cosmologies

Alternatively to the previous scenario, a relaxation to the sphaleron bound can be obtained by modifying the underlying cosmology. In order to do so, we will introduce ourselves into the braneworld language where we live in a brane embedded in a higher dimensional Universe. Within this scenario, one can consider different forms of the Stress-Energy momentum on the Bulk, which lead to the non standard behaviour of the Universe on the brane we are looking for, by suitable choices of boundary conditions.

Regarding braneworlds, Randall-Sundrum argued that an ADS bulk and a brane with negative tension can provide a simple solution to the hierarchy problem [39]. Moreover, the Hubble rate on the brane under this scenario shows a non standard form

$$H^2 = \frac{8\pi}{3M_{Pl}^2}\rho\left(1+\frac{\rho}{2\sigma}\right)+\frac{\mathcal{C}}{a^4}\ . \tag{6.45}$$

On the other hand, Chung and Freese [40] showed that, in the context of braneworlds, it is possible to find any function of the FRW equation if one changes the stress energy tensor composition in the bulk. Therefore, in general, one can parametrize the expansion rate in the following form

$$H^2 = \kappa\rho + \mu\rho^n \tag{6.46}$$

where $\kappa = \frac{8\pi}{3M_{Pl}^2}$ and $\mu \sim \mathcal{O}(\mathrm{GeV}^{-(4n-2)})$. Notice that the geometry of such a Universe is flat and it is trivial to see that each value of n will lead to a different class of FRW equations.

For $n < 2/3$, we find the so-called "Cardassian models" [41], where one can explain the acceleration of a flat Universe at late times. In this work, however, we are interested in the opposite regime for n. We will show that, any n, with $n > 1$ can play an important role reopening the window for electroweak baryogensis without enlarging the particle content. In [42], a study was done for a Randall-Sundrumm like Universe, which in the particular case of $n = 2$. We will generalize this analysis for a generic modified expansion rate, showing that the Randall-Sundrum is only one particular choice among all the possible cases.

For simplicity, we will consider a radiation dominated Universe. There the expansion rate can be written as

$$H^2 = \kappa\rho_r\left(1+\frac{\rho_r^{n-1}}{M^{4(n-1)}}\right)\ , \tag{6.47}$$

where ρ_r is the radiation energy density and M is the scale at which the transition to the usual FRW equation takes place. As explained before, contrary to the Cardassian models, we are seeking for departures of the standard expansion rate at early times and therefore, we need to explore $n > 1$. Remember that $n < 1$, provides late time accelerated expansion and while it gives a nice explanation for a flat, expanding matter dominated universe, it cannot play any role during the electroweak phase transition ($n = 1$ recovers the usual FRW).

The above expression may be rewritten in a straightforward manner as

$$\begin{aligned} H^2 = & \quad \kappa\rho_r(T)\left(1 + \frac{T^{4(n-1)}}{T_m^{4(n-1)}}\right) \\ = & \quad \kappa\rho_r(T)\left[1 + \left(\frac{T}{T_m}\right)^{4(n-1)}\right] , \end{aligned} \tag{6.48}$$

where $\rho_r(T) = \frac{\pi^2}{30} g_* T^4$ and T_m is the matching temperature, the temperature at which we evolve from a Universe with a modified FRW constraint to the usual one.

As we can see, at earlier epochs the second term dominates over the former one. Thus, using the value of the Hubble rate needed to make the phase transition strongly first order, one can find a correlation between the matching temperature and the power of the cardassian model

$$H(T) = H_0(T)\left(\frac{T}{T_m}\right)^{2(n-1)} , \tag{6.49}$$

where $H_0(T) = 1.66 g_*^{1/2} \frac{T^2}{M_{Pl}}$.

Using this expression for the Hubble rate at the electroweak scale, the sphaleron bound can change significantly for different values of T_m and n. This is plotted in figure 6.8.

5 Conclusions

In this work we have shown that despite the fact that the available region in parameter space for the SM and most of its extensions (most notably the MSSM) for electroweak baryogenesis is highly constrained

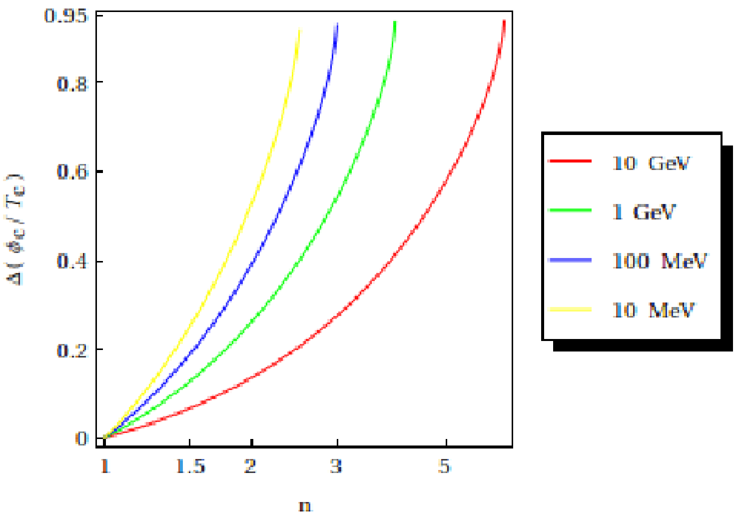

Fig. 6.8: Relative change on the sphaleron bound as a function of n for different matching temperatures to the usual FRW scenario T_m

by experimental results, it can be increased in some alternative scenarios, without enlarging its particle content.

In particular we have discussed a scenario with an early period of matter domination triggered by a long lived massive particle. In such a case, the expansion rate can be orders of magnitude larger than the standard, radiation dominated one and substantially relax the sphaleron bound. The decay of this massive particle generates a huge entropy production that can dilute away any unwanted relic, turning the constraints on the inflationary reheating temperature unnecessary. In this respect we have shown that an early period of matter domination mimics the nice effects of thermal inflation with no additional particle content. However this entropy production also imposses strong constaints on the efficiency of the mechanism for baryogenesis at the electroweak scale.

On the other hand, when analyzing thermal histories suffering very prolonged periods of matter domination preceding the usual one, one wonders whether there is any signature of their existence left that can be tested today. An obvious place to look is of course, structure formation.

The total perturbation amplitude growth during the first matter-dominated phase will just be $a_{\rm end}/a_{\rm start}$. Then, if the primordial perturbation amplitude (say, from inflation) is larger than 10^{-14}, this just means that the structure becomes strongly non- linear for very prolonged periods of matter domination.

However we should keep in mind that the perturbation growth only occurs for perturbations inside the Hubble radius, with the maximum growth occurring for those scales that came inside the Hubble radius at the beginning of the matter dominated epoch, *i.e.* the smallest scales. Scales that entered the horizon later than a time t_i will grow by only $a_{\rm end}/a_i$, where a_i is the scale factor at which they entered the horizon. Those will still be small physical scales today. Scales that never crossed inside the horizon during the early matter dominated epoch would not have this enhancement. So the prediction of this model for the perturbation power spectrum would be the ordinary LCDM + inflation spectrum on large scales with an enhancement of power that grows as a power of wavenumber$k$ on small scales. The enhancement would set in gradually for$k > k_e$, where $k_e = (aH)_{\rm end}$, the comoving wavenumber above which the power spectrum is enhanced. Roughly, in the usual CDM model, the mass power spectrum $P \propto k^{n-4}$ on small scales, where $n = 0.96$ is the primordial spectral index from inflation. In this model with massive particle decay and an early matter dominated epoch, the power spectrum on scales $k > k_e$ will instead go as $P \propto (k/k_e)^n$, *i.e.* the power grows on small scales and becomes non-linear on scales $k > 10^{2.5}k_e$.

However, the massive particle decays. And as we need the universe to become radiation dominated before BBN, the decay products should be relativistic. Relativistic particles will free-stream out of the mini-halos even if they are strongly non-linear in density contrast (their gravitational potentials are still weak). So the structures formed will eventually evaporate, leaving no trace of their existence behind. Of course in scenarios more sophisticated than this simple one, there will be traces of this first period of matter domination left. We will carry out this study elsewhere.

We have also shown that a modification of the FRW equation can lead to expansion rates at early times large enough to relax the sphaleron bound to level consistent with current experimental bounds. In such scenarios the transition to the standard cosmology takes place after the electroweak phase transition and before BBN, not affecting then either structure formation or the age of the universe.

7. STRUCTURE FORMATION DURING AN EARLY PERIOD OF MATTER DOMINATION

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abstract

In this work we show that modifying the thermal history of the Universe by including an early period of matter domination can lead to the formation of astronomical objects. However, the survival of these objects can only be possible if the dominating matter decays to a daughter particle which is not only almost degenerate with the parent particle but also has an open annihilation channel. This requirement translates in an upper bound for the coupling of such a channel and makes the early structure formation viable.

1 Introduction

The Universe known and observed nowadays is a consequence of a long process where the primordial seeds were amplified due to Inflation, a stage of the Universe where its size grew exponentially and left all the observable scales out of the horizon. At re-entry after Inflation termination, the seeds of these scales began to accrete matter and formed the observable astrophysical structures. One of the advantages of the standard cosmological scenario is that it is capable of addressing the whole process since the beginning until the late formation of complex structures.

The first seeds are widely supposed to be quantum fluctuations, amplified during the primordial inflationary era. Working in the fourier space, the component responsible for conducting the early exponential expansion of the Universe develops an inhomogeneous perturbation with a certain length and amplitude that gets frozen when the horizon scale becomes smaller than this length. Such a perturbation is then transmitted to the other components of the Universy by gravity. Furthermore, the amplitude of this perturbation is nearly the same for every component of the Universe, once one assumes that perturbations are adiabatic, as experiments seem to confirm.

Inflation is commonly assumed to be followed by a radiation era once pressure becomes important against gravity. Due to this effect, seeds are unable to attract matter and therefore form structures. As a consequence, the gravitational potential for scales entering the horizon throughout this radiation dominated epoch vanishes, a feature exhibited by the power spectrum for such scales getting suppressed by a factor $\frac{1}{k^3}$.

Since the energy density of radiation is diluted with the expansion of the Universe more rapidly than the energy density of matter, pressure becomes insignificant after the matter-radiation equality point and structures can be formed by matter accretion. At this matter dominated stage, perturbations that enter the horizon start to grow linearly with the scale, attracting more matter until they become non linear and collapse into the observed structures. This effect can be seen in the power spectrum profile of the perturbations, that grows as k. One

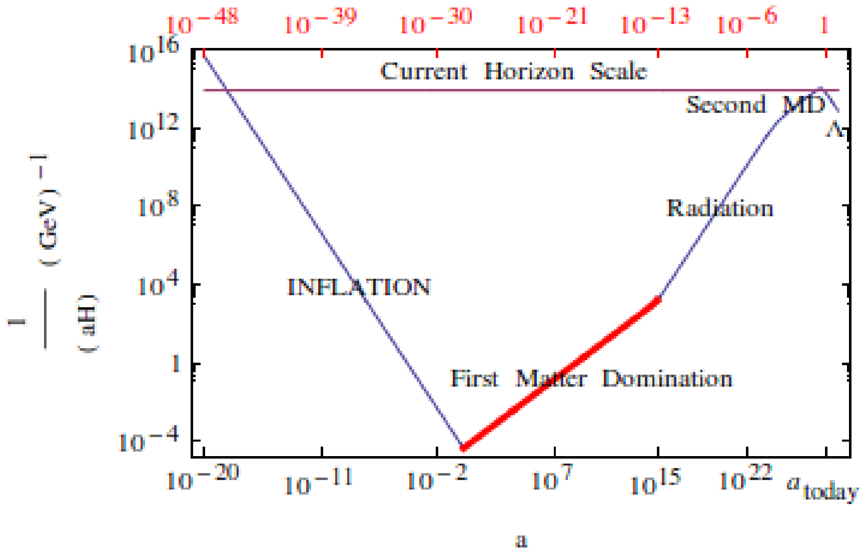

Fig. 7.1: Evolution of the Hubble horizon in a non standard history of the Universe as a function of the scale factor. Scales factors in black (bottom of the plot) correspond to the convention used in this work where $a = 1$ signals the beginning of matter domination. Scale factors in red (top of the plot) correspond to the convention where $a = 1$ is set to today. This double labelling can be used as a "dictionary" for the following plots.

peculiarity of this whole process is that as a consequence of the coupling between baryons and photons, observable matter starts to fall into the gravitational wells at $z_{\rm dec} \approx 1100$, much later than dark matter which, as being weakly interacting, starts to grow and form structures right after the matter-radiation equality time $z_{\rm eq} \approx 3400$. This is why first and older objects are searched in the form of halos or mini-halos of dark matter.

Needless to say, this is the cartoon picture of structure formation assuming the standard thermal history of the Universe. However, as far as the Universe is radiation dominated by BBN, $T_{BBN} \simeq 1$ MeV, and matter-radiation equality takes places at $T_{\rm eq} \simeq 1$ eV, one is free to modify the thermal history at will. For instance, thermal inflation [9] introduces a very short inflationary epoch to get rid of unwanted particles such as moduli. Another case would be to consider that a very heavy particle, with a large thermal abundance, came to dominate the energy density of the Universe at its early stages. This early domination can be healthy for erasing unwanted particles and relaxing the conditions for producing the baryon asymmetry at the electroweak scale [8, 43]. An example of this latter modification in the standard history of the Universe can be seen in figure 7.1, that shows the evolution of the comoving Hubble radius through different eras. Within this sort of picture, an early matter domination era (coloured in red) driven by a heavy particle commences shortly after inflation ends, contrary to the standard picture where inflation and dark matter domination epochs are connected by a long period of radiation domination. One of the most remarkable changes when including such a modification is that scales entering the horizon during this new era can now grow linearly with the scale until their amplitudes become non linear and begin to form substructures, which in principle are to survive up to now. In this work, we study the conditions under which such a scenario can be realised and explore the consequences of such an early structure formation period in a thermal history of the Universe such as the one depicted in figure 7.1.

Our paper will be organised as follows: In section 2, we introduce the setup of the Universe, *i.e*, its components, interactions and magnitudes. Once the setup for such an Universe is given, the history that follows is automatically known. Furthermore, the details of the construction and motivation for such a scenario are also introduced. In section 3, the features of the structure formation picture are explained. We finally conclude in section 4.

2 Scenario details

In [44], a multifluid Universe where a heavy matter particle dominates the thermal history of the Universe until it decays away into radiation and matter was considered. The equation of motions for this case are

$$\frac{\mathrm{d}\rho_{\mathrm{mm}}}{\mathrm{d}t} + 3H\rho_{\mathrm{mm}} = -\Gamma_{\mathrm{mm}}\rho_{\mathrm{mm}} , \quad (7.1)$$

$$\frac{\mathrm{d}\rho_{r}}{\mathrm{d}t} + 4H\rho_{r} = (1 - f_b)\Gamma_{\mathrm{mm}}\rho_{\mathrm{mm}} , \quad (7.2)$$

$$\frac{\mathrm{d}\rho_{dm}}{\mathrm{d}t} + 3H\rho_{dm} = f_b\Gamma_{\mathrm{mm}}\rho_{\mathrm{mm}} , \quad (7.3)$$

$$H^2 = \frac{8\pi}{3M_{Pl}^2}\left(\rho_{mm} + \rho_r + \rho_{dm}\right) , \quad (7.4)$$

where the subscript "mm" stands for mother matter, the component responsible for the early period of matter domination of the Universe, "dm" for daughter matter and f_b is the fraction of mother matter decaying into daughter matter.

Within this kind of Universe, one can reconstruct the history of the Universe by taking suitable values for both Γ_{mm} and f_b. These two parameters are not independent of each other but can be related as follows

$$f_b \simeq \frac{T_{\mathrm{eq}}}{T_{\mathrm{RH}}} , \quad (7.5)$$

$$T_{\mathrm{RH}} \simeq 0.55 g_*^{-1/4}\sqrt{M_{Pl}\Gamma_{\mathrm{mm}}} , \quad (7.6)$$

where T_{RH} is the reheating temperature, *i.e.*, the temperature at which the mother particle releases all its energy and T_{eq} is the temperature when the energy density of radiation and matter are equal, with $T_{\mathrm{eq}} \simeq 1$ eV to get the right amount of dark matter today.

Therefore, if we require that the mother particle has completely decayed away prior to BBN, one may obtain a lower bound on $\Gamma_{\mathrm{mm}} \gtrsim 2.0 \times 10^{-24}$ GeV or likewise an upper bound on $f_b \lesssim 10^{-6}$. Such small values of the branching fraction might be dangerous for the formation of mini-halos. As it was demonstrated by Cen [45], the density of mini-halos $\rho_{\mathrm{mini-halo}}$ decreases by a factor $(f_b)^4$ when a sizeable portion of the main component of such substructures decays into radiation

$$\rho_{\mathrm{mini-halo}}|_f = (f_b)^4 \rho_{\mathrm{mini-halo}}|_i , \quad (7.7)$$

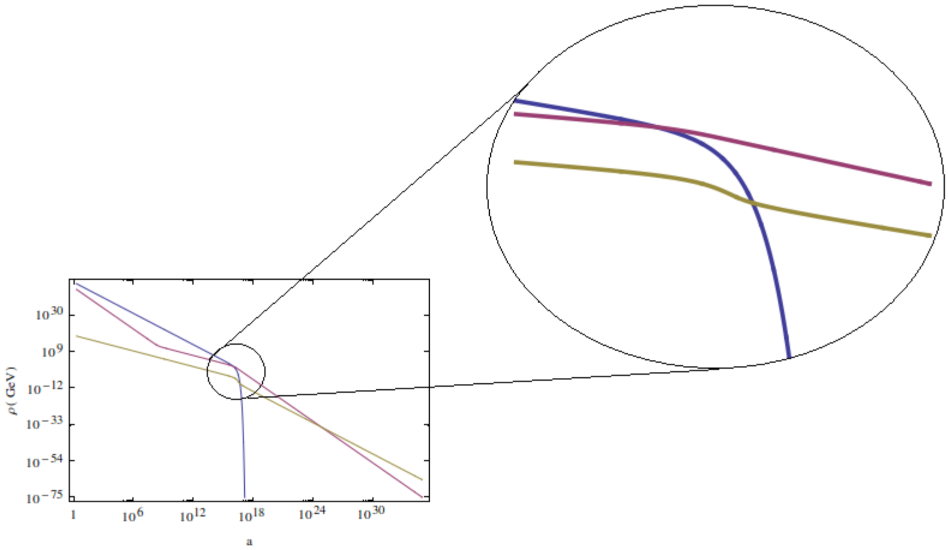

Fig. 7.2: Evolution of the energy densities for the different components of our Universe. The transition between the early matter domination epoch and the standard radiation one is depicted amplified

where the subscript "i" referes to values before the decay and "f" after the decay is completed. Hence, for the mini-halos not to puff up by a large factor, the branching ratio f_b to non-relativistic particles (*i.e.* to the daughter particles) cannot be very small. Consequently, one would need f_b to be not far from unity in order for the substructure to survive. However, as in this scenario the daughter particle is the only dark matter component, then such large values for the branching fraction f_b are forbidden by current observations, as they would lead to an overabundance of dark matter today. Therefore we need to add new channels of entropy production.

Owing to the mentioned argument, our setup will be the same as exposed earlier but including now a channel of annihilation for the daughter matter into radiation. The equations of motion for the energy densities can be then written as follows

$$\begin{aligned}
\frac{\mathrm{d}\rho_{\mathrm{mm}}}{\mathrm{d}t} + 3H\rho_{\mathrm{mm}} &= -\Gamma_{\mathrm{mm}}\rho_{\mathrm{mm}} \ , && (7.8)\\
\frac{\mathrm{d}\rho_{r}}{\mathrm{d}t} + 4H\rho_{r} &= (1-f_b)\Gamma_{\mathrm{mm}}\rho_{\mathrm{mm}} + \Upsilon^{anh} \ , && (7.9)\\
\frac{\mathrm{d}\rho_{dm}}{\mathrm{d}t} + 3H\rho_{dm} &= f_b\Gamma_{\mathrm{mm}}\rho_{\mathrm{mm}} - \Upsilon^{anh} \ , && (7.10)\\
H^2 &= \frac{8\pi}{3M_{Pl}^2}\left(\rho_{mm} + \rho_r + \rho_{dm}\right) \ , && (7.11)
\end{aligned}$$

where

$$\Upsilon^{anh}(t) = \gamma\left(\rho_{dm}^2(t) - \rho_{\mathrm{eq}}^2(t)\right)$$

is the operator for the annihilation of daughter matter into radiation. $\rho_{eq}(t)$ is the equilibrium density. Given the fact that our daughter matter density is not a thermal relic, we have set this density to 0.

The motivation for including such a term is clear. As it was showed before, mini-halos formation throughout a period of matter domination decaying into radiation are suppressed by the fourth power of f_b. Therefore, any substructure formed would be erased given the upper bound from eq.(7.5) unless we include this new channel that alleviates this effect, letting more production of daughter matter and allowing us to have f_b as large as needed. Particularly, one might take f_b equal to 1, a situation in which the annihilation would be the only source of all the radiation in the Universe[1].

In addition to the conditions explained above, one must not only care about the fraction of matter produced during this period, but also about the velocity at which they are expelled after being formed since perturbations might be washed out by great velocities through free streaming. For such an analisis we need to define a scale λ_{fs} [46]

$$\lambda_{fs}(t) = \int_{t_{reh}}^{t} \frac{< v_{dm} >}{a}\, dt \qquad (7.12)$$

[1] This particular case however would require a reheating temperature of several hundred GeVs or higher. As we will see, as the size of the objects formed during this early matter domination era obviosuly depends on the length of this era, this extremal case with $f_b \sim 1$ is clearly not favoured

below which perturbations get erased away.

If one assumes that the velocity after reheating is diluted linearly by the expansion of the Universe and that the free streaming scales barely change after matter-radiation equality, one finds that [47]

$$\begin{aligned}\lambda_{fs}(a) &= \frac{< v_{RH} > a_{RH}}{H_0\sqrt{\Omega_{rad}}}\int_{a_{reh}}^{a}\frac{1}{a'\sqrt{1+a_{eq}/a'}}\,da' \\ &= \frac{2 < v_{RH} > a_{RH}}{H_0\sqrt{\Omega_{rad}}}\left(\text{arcsinh}\sqrt{\frac{a_{eq}}{a_{RH}}} - \text{arcsinh}\sqrt{\frac{a_{eq}}{a}}\right)\;,\end{aligned} \tag{7.13}$$

where $< v_{RH} > \equiv < v_{dm}(a_{RH}) >$ is the average velocity of the daughter particle at the reheating moment (within the instantaneous decay approximation), H_0 is the current Hubble constant and Ω_{rad} the observed current abundance of radiation.

Regarding the velocity of the daughter particle and assuming that one mother produces a pair of daughters, it can be then easily demonstrated by kinematics that

$$v_{dm}^2 = \left(1 - \frac{4m_{dm}^2}{M^2}\right)\;, \tag{7.14}$$

where v_{dm} is given in units of c, M is the mass of the mother particle and m_{dm} the mass of the daughter matter. Consequently, one can see that in order for the daughter particle to have low velocities when created, so that to avoid free streaming washout effects, it needs to be nearly half of the mass of the mother particle.

On the other hand, we still need the daughter particle to give rise to the right amount of observed dark matter. Once the mother matter decays away completely at $T_{\rm RH}$, the density of daughter matter can be written as $\rho_{\rm dm} \simeq f_b\rho_{\rm rad}$. At this point, the annihilation term dominates over the expansion term in the equation for $\rho_{\rm dm}$, making it decay abruptly until both terms balance. This effect takes place shortly after the reheating time, where $\rho_{\rm dm} \simeq \frac{H(a_{\rm RH})}{\gamma}$. From that point onwards, the remaining density dilutes in the standard way with the expansion of the Universe to provide the observed amount of matter.

This allows us to constrain the size of the annihilation coupling to be

$$\gamma \simeq 5\times 10^{-1}\frac{1}{(M_{Pl}\;T_{\rm RH}\;T_{\rm eq})} = 5,5\times 10^{-8}\left(\frac{1\ {\rm MeV}}{T_{\rm RH}}\right)\quad {\rm GeV}^{-3}\ . \quad (7.15)$$

In summary, the conditions/ingredients any model leaving traces of an earlier epoch of matter domination should have, are the following

1. A heavy particle, that we have called "mother", dominates the energy density of the Universe up to its decay into radiation and matter. The latter one is labeled as "daughter".
2. The daughter particle would be the candidate for WIMP dark matter.
3. The heavy mother particle forms mini-halos during the first matter dominated era before it decays. In order for those structures not to evaporate completely, the daughter particles must be borned non- relativistic, *i.e.*, their masses must be nearly degenerate with that of the mother ($m \approx \frac{M}{2}$).
4. For the mini-halos not to puff up by a large factor, the branching ratio f_b to non-relativistic particles (*i.e.*, to the daughter particle) cannot be very small since the density of the mini-halo decreases by a factor $(f_b)^4$.
5. On the other hand, in order to have a radiation dominated universe by the time of BBN and until the usual epoch of matter domination, f_b cannot be extremely large if it is the only source of entropy production.
6. It is impossible to simultaneously satisfy the last two conditions unless one includes a annihilation term, whose size is constrained by equation 7.15.

3 Perturbations and structure formation

From the previous section, we have seen that it is plausible to have an Universe dominated by a very heavy particle which finally decayed

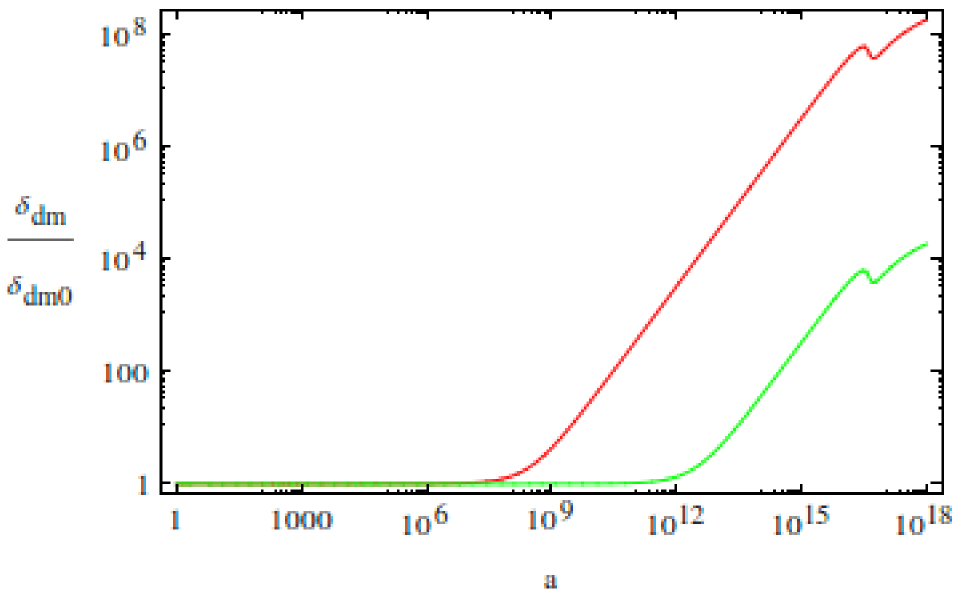

Fig. 7.3: Evolution in scale factor of the density contrast of the daughter particle, our would be dark matter candidate, for two different scales in units of the initial perturbation. Red line corresponds to $k = 10^4 k_{RH}$ and green line to $k = 100 k_{RH}$. Structures become non linear for $(\delta_{dm}/\delta_{dm0}) \sim 10^5$ corresponding to $\delta_{dm0} \sim 10^{-5}$ as seen by CMB measurements. The arbitrary initial value for the scale factor has been taken equal to 1 when solving the equations of motion.

into radiation and common matter. Moreover, as it was also pointed out, density perturbations entering the horizon during that epoch, can grow significantly until the non-linear regime is reached and form substructures. These substructures are very sensitive to the production of entropy, so high abundances of radiation during this epoch may delete any substructure formed. Therefore, we added an annihilation term for the daughter matter, which will mainly act after the heavy particle decayed away allowing to have less amount of radiation during this early matter domination epoch.

The equations for the density perturbations read as follows

$$a^2E(a)\delta'_{\rm mm}(a)+\tilde\theta_{\rm mm}(a)+3a^2E(a)\Phi'(a)=a\tilde\Gamma_{\rm mm}\Phi(a)\ , \tag{7.16a}$$

$$a^2E(a)\tilde\theta'_{\rm mm}(a)+aE(a)\tilde\theta_{\rm mm}+\tilde k^2\Phi(a)=0\ , \tag{7.16b}$$

$$\begin{aligned}a^2E(a)\delta'_r(a)+\frac{4}{3}\tilde\theta_r(a)+4a^2E(a)\Phi'(a)&=(1-f)\frac{\tilde\rho^0_{\rm mm}(a)}{\tilde\rho^0_r(a)}a\tilde\Gamma_{\rm mm}\left[\delta_{\rm mm}(a)-\delta_r(a)-\Phi(a)\right]+\\&+\frac{a}{H_1}\frac{\gamma}{\rho_r}\left[\left(\rho^2_{dm}-\rho^2_{eq}\right)(\delta_r+\Phi)-2\delta_{dm}\rho^2_{dm}\right]\ ,\end{aligned} \tag{7.16c}$$

$$\begin{aligned}a^2E(a)\tilde\theta'_r(a)+\tilde k^2\Phi(a)-\tilde k^2\frac{\delta_r(a)}{4}&=(1-f)\frac{\tilde\rho^0_{\rm mm}(a)}{\tilde\rho^0_r(a)}a\tilde\Gamma_{\rm mm}\left[\frac{3}{4}\tilde\theta_{\rm mm}(a)-\theta_r(a)\right]+\\&+\frac{a}{H_1}\frac{\gamma\left(\rho^2_{dm}-\rho^2_{eq}\right)}{\rho^0_r}\left[-\frac{3}{4}\theta_{dm}+\theta_r\right]\ ,\end{aligned} \tag{7.16d}$$

$$\begin{aligned}a^2E(a)\delta'_{\rm dm}(a)+\tilde\theta_{\rm dm}(a)+3a^2E(a)\Phi'(a)&=f\frac{\tilde\rho^0_{\rm mm}(a)}{\tilde\rho^0_{\rm dm}(a)}a\tilde\Gamma_{\rm mm}\left[\delta_{\rm mm}(a)-\delta_{\rm dm}(a)-\Phi(a)\right]+\\&+\frac{a}{H_1}\left(-\frac{\gamma}{\rho_{dm}}\right)\left[\left(\rho^2_{dm}-\rho^2_{eq}\right)(\delta_{dm}+\Phi)-2\delta_{dm}\rho^2_{dm}\right]\ ,\end{aligned} \tag{7.16e}$$

$$a^2E(a)\tilde\theta'_{\rm dm}(a)+aE(a)\tilde\theta_{\rm dm}+\tilde k^2\Phi(a)=f\frac{\tilde\rho^0_{\rm mm}(a)}{\tilde\rho^0_{\rm dm}(a)}a\tilde\Gamma_{\rm mm}\left[\tilde\theta_{\rm mm}(a)-\tilde\theta_{\rm dm}(a)\right]\ , \tag{7.16f}$$

$$\tilde k^2\Phi+3aE^2(a)\left[a^2\Phi'(a)+a\Phi(a)\right]=\frac{3}{2}a^2\left[\tilde\rho^0_{\rm mm}(a)\delta_{\rm mm}(a)+\tilde\rho^0_r(a)\delta_r(a)+\tilde\rho^0_{\rm dm}\delta_{\rm dm}(a)\right]\ , \tag{7.16g}$$

where $E(a)\equiv\frac{H(a)}{H_0}$, $\tilde k\equiv\frac{k}{H_0}$, $\tilde\theta_{\{mm,dm,r\}}\equiv\frac{\theta_{\{mm,dm,r\}}}{H_0}$ and $\tilde\rho_{\{mm,dm,r\}}\equiv\frac{\rho_{\{mm,dm,r\}}}{\rho_0}$ with H_0 and ρ_0 being the initial Hubble rate and total energy density of the Universe respectively. The details about the derivation of these equations are given in appendix 5.2.

The equations given above reproduce the ones in [44], once the annihilation terms are set to zero. It can be seen that these terms source the equations for the density perturbations and velocities of the daughter and radiation component respectively (eqs. (7.16c-7.16f)), playing a fundamental role when the mother component is on the verge of decaying and allowing the daughter particle to release part of its energy into radiation. This entropy production can be seen in figure 7.3, where perturbations in the energy density of the daughter particle, *i.e.* what would be our dark matter candidate, are depicted and we can see that they tend to decrease several orders of magnitude when the mother particle decays away completely. Fortunately, as it was already mentioned, this decrease takes place near the reheating point, so it is expected that any density perturbations which have already entered in the non-linear regime will survive although their size can slightly decrease. In addition, one should notice that in both figures, what is

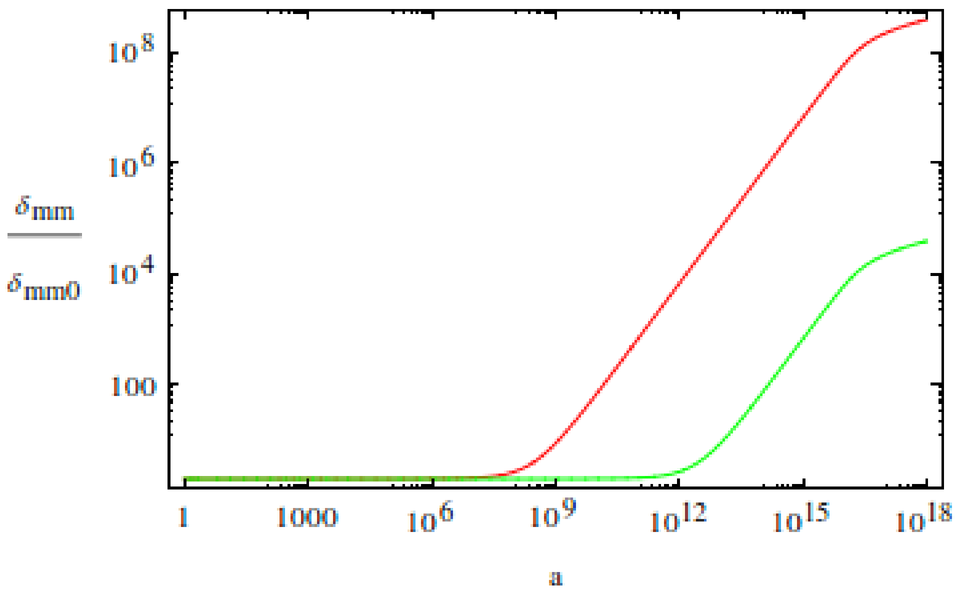

Fig. 7.4: Evolution in scale factor of the density contrast of the mother particle for two different scales in units of the initial perturbation. Red line corresponds to $k = 10^4 k_{RH}$ and green line to $k = 100k_{RH}$. Structures become non linear for $(\delta_{mm}/\delta_{mm0}) \sim 10^5$ corresponding to $\delta_{mm0} \sim 10^{-5}$ as seen by CMB measurements. The arbitrary initial value for the scale factor has been taken equal to 1 when solving the equations of motion.

plotted is the density contrast, *i.e.* the size of the perturbations in terms of its initial size. Such a quantity is determined by Inflation but anisotropy measurements in the CMB map set it to be $\delta_{dm0} \approx 10^{-5}$, which in the standard case of adiabatic perturbations, takes this value for all the components of the Universe[2]. Therefore, the non-linear regime would be reached in our figures when δ_{dm} times the size of the seeds is of order one.

Likewise, it is important to clarify that even the smallest scales which enter the horizon during this early matter domination epoch and

[2] We are assuming that the nearly scale invariance of primordial perturbations from Inflation still holds for such small scales

do not have enough time to reach the non-linear regime (the example in figure 7.3 with $k = 100k_{RH}$), **will still experience a remarkable growth, which will lead to the formation of structures and substructures much earlier than in the standard picture once the Universe becomes matter dominated again.** In particular, this feature may be important for a Dark Matter Halo to collapse shortly after matter-radiation equality forming an ultracompact minihalo, which are excellent indirect detection targets [48, 49] and attractive for lensing prospects [50].

In figure 7.4 the evolution of the density perturbations of the mother particle is shown. As it can be seen, such an evolution, behaving as matter, is very similar to that of the daughter except for the effect coming from the annihilation channel, which is absent in this component.

Finally, the evolution of the radiation perturbations are plotted in figure 7.5. As we can see, the amplitude is amplified during the early matter domination epoch until it decreases completely and begins to oscillate with a negligible value when the mother particle releases all the energy. Such behaviour is very similar to the one given in [44], a fact which exhibits that the annihilation channel has a minor effect on radiation perturbations.

Regarding the density perturbations and structure formation, a variable to study is σ, the variance of the density perturbations smoothed at a certain scale, normally used to analise, within the Press-Schechter formalism [51], the abundance and evolution of halos and sub-halos at relevant scales and with a certain size, which is given by

$$\sigma^2_{\mathrm{RH}}(R) = \int_0^\infty \frac{dk}{k} \left(\frac{k}{a_{\mathrm{RH}} H(a_{\mathrm{RH}})} \right)^4 W^2(kR) T^2(k) \delta^2_H(k) \ , \qquad (7.17)$$

where the subscript RH means that this quantity is evaluated at the reheating time when the mother particle releases all the energy.

Let us briefly explain the formula (7.17). As it was already mentioned, σ^2 is the density perturbations smoothed for a certain scale R. This role is played by the function $W(kR)$, which is responsible for filtering out those modes with $kR \geq 1$ and therefore allow us to study

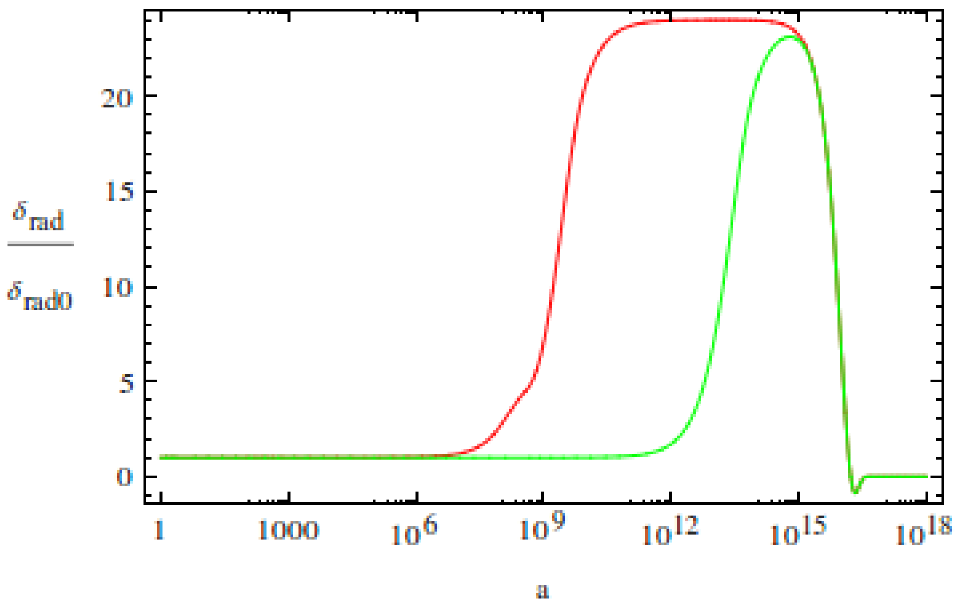

Fig. 7.5: Evolution in scale factor of the density constrast of radiation for two different scales. Red line corresponds to $k = 10^4 k_{RH}$ and green line to $k = 100k_{RH}$. The initial contrast value corresponds to $\delta_{rad0} \simeq 10^{-5}$ as seen by CMB experiments. The arbitrary initial value for the scale factor has been taken equal to 1 when solving the equations of motion.

the relevant scales. In order to do this, we have used the following filter function

$$W(kR) = \exp(-\frac{1}{2}k^2(\alpha R)^2) \times W_{\text{top-hat}}(kR) \; , \qquad (7.18)$$

where $W_{\text{top-hat}}(kR) = \frac{3}{(kR)^3}\left[\sin(kR) - (kR)\cos(kR)\right]$ is the usual top-hat window function. For our purposes, however, we wish to focus upon the scales that enter the horizon during the early matter domination epoch and this is not achieved with the usual top-hat window function. Owing to this, we introduced an exponential function to suppress modes with $k < k_{RH}$.

On the other hand, $T(k)$ is the well known transfer function which

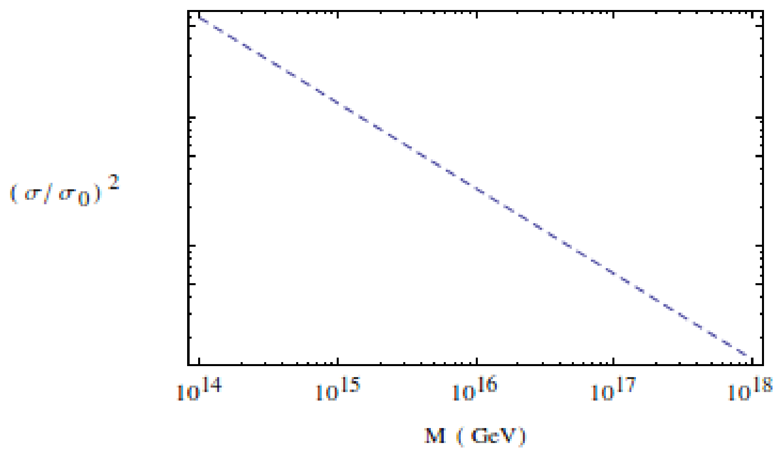

Fig. 7.6: The variance of the daughter density perturbations σ in terms of the mass M contained in a sphere R. σ_0 is the normalization factor accounting for the several constants appearing in the equation 7.17. The scale in both axes is logarithmic.

for the scales that we are taking into account, is scale invariant [44] and $\delta_h(k)$ is the amplitude of the primordial density perturbations originated during inflation, which can be written as

$$\delta_h(k) = 1.87 \times 10^{-5} \left(\frac{k}{k_{pivot}}\right)^{\frac{(n_s-1)}{2}} , \tag{7.19}$$

where $k_{pivot} = 0.002$ Mpc^{-1} and n_s is the spectral index $n_s = 0.9603 \pm 0.0073$ [11].

Finally, the factor $\left(\frac{k}{a_{\rm RH} H(a_{\rm RH})}\right)^4$ takes into account the scale factor growth of modes entering during matter domination.

In figure 7.6, we show the normalised σ^2 (in arbitrary units) evaluated at the reheating epoch for different mass objects, related to their size by $\rho = \frac{M}{\frac{4\pi}{3}R^3}$, where ρ is the total energy density at that moment. As it can be seen, this quantity is a mass decreasing function, meaning that the population of heavier objects is lower since they correspond to scales that entered later in the horizon and thus had less time to become non linear and begin to accrete matter.

The size of the dark matter objects formed is very model dependent, however to get a flavour of it, it is worth remembering that the equivalent horizon mass scale at the QCD epoch ($T \sim 100$ MeV) is around the mass of Jupiter. Moreover, one can easily work out the comoving Hubble size at the reheating time in terms of current parameters as

$$k_{RH}^{-1} \sim 10^{-6}\sqrt{\frac{\Omega_r}{\Omega_m}}\left(\frac{1\ \mathrm{MeV}}{T_{RH}}\right)k_0^{-1} \, . \tag{7.20}$$

Plugging the today known parameters, one can find that if the first matter domination era ends before BBN, a scale size which corresponds to roughly a parsec, then the significant power enhancement (*i.e.* formation of non-linear structure with perturbation amplitude of unity during the first matter domination era) would be on somewhat smaller scales than that, presumably corresponding to the milliparsec regime or even planets and stellar masses.

4 Conclusions

In this work, we have shown that the formation of observational objects can be very sensitive to changes in the thermal history of the Universe. In particular, an early period of matter domination could amplify the primordial inflationary seeds leading to the formation of halos or mini-halos, objects which can be in principle observable and detectable [48, 49, 52–54].

Since at some point by BBN one needs to recover the usual picture of a radiation dominated Universe, one needs to care about the transition between both phases due to the production of entropy. Such production may erase or at least reduce any structure formed during the early period of matter. In particular, the density of primordial objects is suppressed by the fourth power of the branching fraction into radiation of the leading component during the matter epoch. As one needs to connect this scenario with the usual picture, *i.e.* radiation domination by BBN and right amount of dark matter abundance, one is forced to using values of the branching ratio which dilutes any primordial objects.

In order to solve this, we have introduced a new channel for the annihilation of the daughter matter into radiation. This allows us to have less amount of radiation during the period of structure formation and thus, larger values of the branching function. Furthermore, as it was showed in the profile of figure 7.2, this new channel only plays an important role when all the energy of the mother particle is totally released, connecting the end of the early matter domination era with the usual picture, and therefore any dilution can only take place when perturbations have already entered in the non-linear regime. We have also showed that this will only happen for modes that entered the horizon early enough to fall into the non-linear regime. In terms of the scale factor, it will happen for $(a_{RH}/a) \gtrsim 10^6$ or $(k/k_{RH}) \gtrsim 10^3$. Modes with $1000k_{RH} > k > k_{RH}$, *i.e.* that don't reach the non-linear regime before the decay of the mother particle, enter the non-linear regime in the second matter domination epoch, but may start to collapse into potential wells much earlier than within the standard thermal history picture due to the earlier growth.

We have also estimated that the new objects beginning to form during this first matter dominated epoch correspond to the milliparsec regime. Can such small scales have any observational relevance for the CMB?. In principle, it is hard to tell since one needs to evolve the perturbations after they entered in the non-linear regime all the way throughout the radiation dominated epoch. Certainly, there are many intriguing features and potentially interesting signatures for models with a (long enough) early period of matter domination able to leave potentially observable substructures. Of course a complete analysis needs to be performed by making use of non-linear methods such as N-Body simulations and falls beyond the scope of this manuscript. Hopefully, our work will trigger such an analysis and above all will let the reader judge himself the grade of apprehension that is appropiate when examining the phenomenology of these theories that take us away from the standard thermal history of the Universe.

To make this picture complete, one may argue that the today existing dark matter abundance does not come primarily from the decay of the daughter particle but from the freeze-out of non relativisitic matter from thermal equilibrium. This would require smaller branching fractions f_b, with the consequent creation of even more entropy which would erase more easily the substructures formed. Moreover, given the features of many proposed candidates for dark matter in the freeze-out scenario (specially neutralino), one would need larger reheating temperatures which would then lead to a less prolonged period of matter domination. As it is at this epoch when perturbations can grow until they enter in the non-linear regime, a dark matter relic density coming only from the decay of a heavy particle appears to be the most favourable scenario regarding an early structure formation in the Universe.

Finally, one may wonder how the observed baryon asymmetry is generated in a scenario like this. At first sight, it seems that it can only come from the decay of the mother particle, imposing more restrictions on its properties. A mechanism viable with having such a heavy particle could be a net baryon number production by means of a derivative coupling of the mother particle to the lepton/baryon current. Such an operator yields an effective chemical potential for baryons and anti-baryones when CPT is violated, allowing the velocity of the heavy particle to develop a non-zero vacuum expectation value [19, 55, 56]. Alternatively, one could also resort to the electroweak phase transition to produce the baryon asymmetry by changing the underlying thermal history of the Universe to being matter dominated during the EWPT, which requires an efficient baryogenesis mechanism due to the entropy production [43].

Acknowledgments

It is a pleasure to thank Scott Dodelson and Will Kinney for useful comments. The authors acknowledge financial support from spanish MEC and FEDER (EC) under grant FPA2011-23596, and Generalitat Valenciana under the grant PROMETEO II/2013/017. GB acknowledges partial support from the European Union FP7 ITN INVISIBLES (Marie Curie Actions, PITN- GA-2011- 289442).

5 *Appendix*

5.1 *Pertubation equations without annhiliation*

In this section, we derive the perturbation equations when an operator for the decaying of the mother particle is added.

In general, the energy conservation equation can be written in a covariant way as follows

$$\nabla_\mu \left({}^{(i)}T^\mu_\nu \right) = Q_\nu \ . \tag{7.21}$$

For the case of decaying matter we have the following

$$Q^{(\phi)}_\nu = {}^{(\phi)}T_{\mu\nu} u^\mu_\phi \Gamma_\phi \ , \tag{7.22}$$

$$Q^{(r)}_\nu = -(1-f) Q^{(\phi)}_\nu \ , \tag{7.23}$$

$$Q^{(dm)}_\nu = -f Q^{(\phi)}_\nu \ , \tag{7.24}$$

where f_b is the branching fraction, Γ_ϕ is the decay operator, $u^\mu_\phi = \left(1-\psi, \vec{V}\right)$ is the perturbed 4-velocity and $T_{\mu\nu}$ is the stress energy-momentum tensor, which in the perfect fluid case reads

$$T^{\mu\nu} = (\rho + P) u^\mu u^\nu + P g^{\mu\nu} \ . \tag{7.25}$$

We shall work in the Newtonian gauge of the perturbed FRW metric, which reads as

$$\mathrm{d}s^2 = -(1+2\Psi)\mathrm{d}t^2 + a^2(t)\delta_{ij}(1+2\Phi)\mathrm{d}x^i \mathrm{d}x^j \ . \tag{7.26}$$

With the above ingredients, one is able to derive the perturbation equations [44]

$$a^2 E(a)\delta'_{\rm mm}(a) + \tilde{\theta}_{\rm mm}(a) + 3a^2 E(a)\Phi'(a) = a\tilde{\Gamma}_{\rm mm}\Phi(a), \tag{7.27}$$

$$a^2 E(a)\tilde{\theta}'_{\rm mm}(a) + aE(a)\tilde{\theta}_{\rm mm} + \tilde{k}^2\Phi(a) = 0, \tag{7.28}$$

$$a^2 E(a)\delta'_r(a) + \frac{4}{3}\tilde{\theta}_r(a) + 4a^2 E(a)\Phi'(a) = (1-f)\frac{\tilde{\rho}^0_{\rm mm}(a)}{\tilde{\rho}^0_r(a)} a\tilde{\Gamma}_{\rm mm}\left[\delta_{\rm mm}(a) - \delta_r(a) - \Phi(a)\right], \tag{7.29}$$

$$a^2 E(a)\tilde{\theta}'_r(a) + \tilde{k}^2\Phi(a) - \tilde{k}^2\frac{\delta_r(a)}{4} = (1-f)\frac{\tilde{\rho}^0_{\rm mm}(a)}{\tilde{\rho}^0_r(a)} a\tilde{\Gamma}_{\rm mm}\left[\frac{3}{4}\tilde{\theta}_{\rm mm}(a) - \theta_r(a)\right], \tag{7.30}$$

$$a^2 E(a)\delta'_{\rm dm}(a) + \tilde{\theta}_{\rm dm}(a) + 3a^2 E(a)\Phi'(a) = f\frac{\tilde{\rho}^0_{\rm mm}(a)}{\tilde{\rho}^0_{\rm dm}(a)} a\tilde{\Gamma}_{\rm mm}\left[\delta_{\rm mm}(a) - \delta_{\rm dm}(a) - \Phi(a)\right], \tag{7.31}$$

$$a^2 E(a)\tilde{\theta}'_{\rm dm}(a) + aE(a)\tilde{\theta}_{\rm dm} + \tilde{k}^2\Phi(a) = f\frac{\tilde{\rho}^0_{\rm mm}(a)}{\tilde{\rho}^0_{\rm dm}(a)} a\tilde{\Gamma}_{\rm mm}\left[\tilde{\theta}_{\rm mm}(a) - \tilde{\theta}_{\rm dm}(a)\right] \ , \tag{7.32}$$

$$\tilde{k}^2\Phi + 3aE^2(a)\left[a^2\Phi'(a) + a\Phi(a)\right] = \frac{3}{2}a^2\left[\tilde{\rho}^0_{\rm mm}(a)\delta_{\rm mm}(a) + \tilde{\rho}^0_r(a)\delta_r(a) + \tilde{\rho}^0_{\rm dm}\delta_{\rm dm}(a)\right], \tag{7.33}$$

where $E(a) \equiv \frac{H(a)}{H_0}$, $\tilde{k} \equiv \frac{k}{H_0}$, $\tilde{\theta}_{\{mm,dm,r\}} \equiv \frac{\theta_{\{mm,dm,r\}}}{H_0}$ and $\tilde{\rho}_{\{mm,dm,r\}} \equiv \frac{\rho_{\{mm,dm,r\}}}{\rho_0}$ with H_0 and ρ_0 being the initial Hubble rate and total energy density of the Universe respectively.

5.2 Pertubation equations with annhiliation

We will now focus on the modification of the density perturbation equations when including an annihilation term.

If we now add a source term accounting for the annhiliation of matter into radiation, these equations would be given as

$$Q_\nu^{(\phi)} =^{(\phi)} T_{\mu\nu} u_\phi^\mu \Gamma_\phi \,, \tag{7.34}$$

$$Q_\nu^{(r)} = -(1-f) Q_\nu^{(\phi)} + Q_\nu^{anh} \,, \tag{7.35}$$

$$Q_\nu^{(dm)} = -f Q_\nu^{(\phi)} - Q_\nu^{anh} \,. \tag{7.36}$$

So our ansatz for the annihilation source could be the following

$$Q_\nu^{anh} = -\gamma \left({}^{(dm)}T_\nu{}^r \; {}^{(dm)}T_{r\mu} - \rho_{eq}^2 \; g_{\nu\mu} \right) u^\mu \,. \tag{7.37}$$

It can be seen that the zero component at zero order gives rise to the right operator

$$\begin{aligned} Q_0^{anh} &= -\gamma \left(g_{0\lambda} \; {}^{(dm)}T^{\lambda r} g_{rs} g_{\mu w} \; {}^{(dm)}T^{sw} - \rho_{eq}^2 \; g_{0\mu} \right) u^\mu \\ &= -\gamma \left(g_{0\lambda} \; \rho^{dm} u^\lambda u^r g_{rs} g_{\mu w} \rho^{dm} u^s u^w - \rho_{eq}^2 \; g_{0\mu} \right) u^\mu \\ &= -\gamma \left(g_{0\lambda} \; \rho^{2,dm} (u_s \cdot u^s) u^\lambda (u_\mu \cdot u^\mu) - \rho_{eq}^2 \; g_{0\mu} u^\mu \right) \\ &= -\gamma \; g_{00} \left(\rho^{2,dm} - \rho_{eq}^2 \right) u^0 \\ &= +\gamma \left(\rho^{2,dm} - \rho_{eq}^2 \right) \,, \end{aligned} \tag{7.38}$$

where the relations $(u_s \cdot u^s) = -1$, $g_{00} = -1$ y $u^0 = 1$ have been used.

A first order in perturbations Q_ν^{inh} takes then the following form

$$Q_0^{inh} = \gamma \left[\left(\rho^{2,dm} - \rho_{eq}^2 \right) (1 + \Psi) + 2\delta^{dm} \rho^{2,dm} \right] \,, \tag{7.39}$$

$$Q_i^{inh} = -a^2 \gamma V_i \left(\rho^{2,dm} - \rho_{eq}^2 \right) \,. \tag{7.40}$$

Working out the energy conservation equation for each component with the perturbed metric in the Newtonian gauge, one can derive the perturbation equations but including now the annihilation terms. These equations read as follows

$$a^2E(a)\delta'_{\rm mm}(a) + \tilde\theta_{\rm mm}(a) + 3a^2E(a)\Phi'(a) = a\tilde\Gamma_{\rm mm}\Phi(a)\ , \tag{7.41}$$

$$a^2E(a)\tilde\theta'_{\rm mm}(a) + aE(a)\tilde\theta_{\rm mm} + \tilde k^2\Phi(a) = 0\ , \tag{7.42}$$

$$a^2E(a)\delta'_r(a) + \frac{4}{3}\tilde\theta_r(a) + 4a^2E(a)\Phi'(a) = \cdots + \frac{a}{H_1}\frac{1}{\rho_r^0}\left[Q_0^{anh,(0)}\delta_r - Q_0^{anh,(1)}\right]\ , \tag{7.43}$$

$$a^2E(a)\tilde\theta'_r(a) + \tilde k^2\Phi(a) - \tilde k^2\frac{\delta_r(a)}{4} = \cdots + \frac{a}{H_1}\frac{1}{\rho_r^0}\left[\frac{\partial_i Q_i^{anh}}{a(1+w_{rad})} + Q_0^{anh,(0)}\theta_r\right]\ , \tag{7.44}$$

$$a^2E(a)\delta'_{\rm dm}(a) + \tilde\theta_{\rm dm}(a) + 3a^2E(a)\Phi'(a) = \cdots - \frac{a}{H_1}\frac{1}{\rho_{dm}^0}\left[Q_0^{anh,(0)}\delta_{dm} - Q_0^{anh,(1)}\right]\ , \tag{7.45}$$

$$a^2E(a)\tilde\theta'_{\rm dm}(a) + aE(a)\tilde\theta_{\rm dm} + \tilde k^2\Phi(a) = \cdots + 0 \tag{7.46}$$

$$\tilde k^2\Phi + 3aE^2(a)\left[a^2\Phi'(a) + a\Phi(a)\right] = \frac{3}{2}a^2\left[\tilde\rho^0_{\rm mm}(a)\delta_{\rm mm}(a) + \tilde\rho^0_r(a)\delta_r(a) + \tilde\rho^0_{\rm dm}\delta_{\rm dm}(a)\right]\ , \tag{7.47}$$

where $(\dots)$ contains the terms without annihilation. If we show them explicitly, the equations of motions are written as follows

$$a^2E(a)\delta'_{\rm mm}(a) + \tilde\theta_{\rm mm}(a) + 3a^2E(a)\Phi'(a) = a\tilde\Gamma_{\rm mm}\Phi(a)\ , \tag{7.48}$$

$$a^2E(a)\tilde\theta'_{\rm mm}(a) + aE(a)\tilde\theta_{\rm mm} + \tilde k^2\Phi(a) = 0\ , \tag{7.49}$$

$$\begin{aligned} a^2E(a)\delta'_r(a) + \frac{4}{3}\tilde\theta_r(a) + 4a^2E(a)\Phi'(a) &= (1-f)\frac{\tilde\rho^0_{\rm mm}(a)}{\tilde\rho^0_r(a)}a\tilde\Gamma_{\rm mm}\left[\delta_{\rm mm}(a) - \delta_r(a) - \Phi(a)\right] + \\ &+ \frac{a}{H_1}\frac{\gamma}{\rho_r}\left[\left(\rho_{dm}^2 - \rho_{eq}^2\right)(\delta_r + \Phi) - 2\delta_{dm}\rho_{dm}^2\right]\ , \end{aligned} \tag{7.50}$$

$$\begin{aligned} a^2E(a)\tilde\theta'_r(a) + \tilde k^2\Phi(a) - \tilde k^2\frac{\delta_r(a)}{4} &= (1-f)\frac{\tilde\rho^0_{\rm mm}(a)}{\tilde\rho^0_r(a)}a\tilde\Gamma_{\rm mm}\left[\frac{3}{4}\tilde\theta_{\rm mm}(a) - \theta_r(a)\right] + \\ &+ \frac{a}{H_1}\frac{\gamma\left(\rho_{dm}^2 - \rho_{eq}^2\right)}{\rho_r^0}\left[-\frac{3}{4}\theta_{dm} + \theta_r\right]\ , \end{aligned} \tag{7.51}$$

$$\begin{aligned} a^2E(a)\delta'_{\rm dm}(a) + \tilde\theta_{\rm dm}(a) + 3a^2E(a)\Phi'(a) &= f\frac{\tilde\rho^0_{\rm mm}(a)}{\tilde\rho^0_{\rm dm}(a)}a\tilde\Gamma_{\rm mm}\left[\delta_{\rm mm}(a) - \delta_{\rm dm}(a) - \Phi(a)\right] + \\ &+ \frac{a}{H_1}\left(-\frac{\gamma}{\rho_{dm}}\right)\left[\left(\rho_{dm}^2 - \rho_{eq}^2\right)(\delta_{dm} + \Phi) - 2\delta_{dm}\rho_{dm}^2\right]\ , \end{aligned} \tag{7.52}$$

$$a^2E(a)\tilde\theta'_{\rm dm}(a) + aE(a)\tilde\theta_{\rm dm} + \tilde k^2\Phi(a) = f\frac{\tilde\rho^0_{\rm mm}(a)}{\tilde\rho^0_{\rm dm}(a)}a\tilde\Gamma_{\rm mm}\left[\tilde\theta_{\rm mm}(a) - \tilde\theta_{\rm dm}(a)\right]\ , \tag{7.53}$$

$$\tilde k^2\Phi + 3aE^2(a)\left[a^2\Phi'(a) + a\Phi(a)\right] = \frac{3}{2}a^2\left[\tilde\rho^0_{\rm mm}(a)\delta_{\rm mm}(a) + \tilde\rho^0_r(a)\delta_r(a) + \tilde\rho^0_{\rm dm}\delta_{\rm dm}(a)\right]\ , \tag{7.54}$$

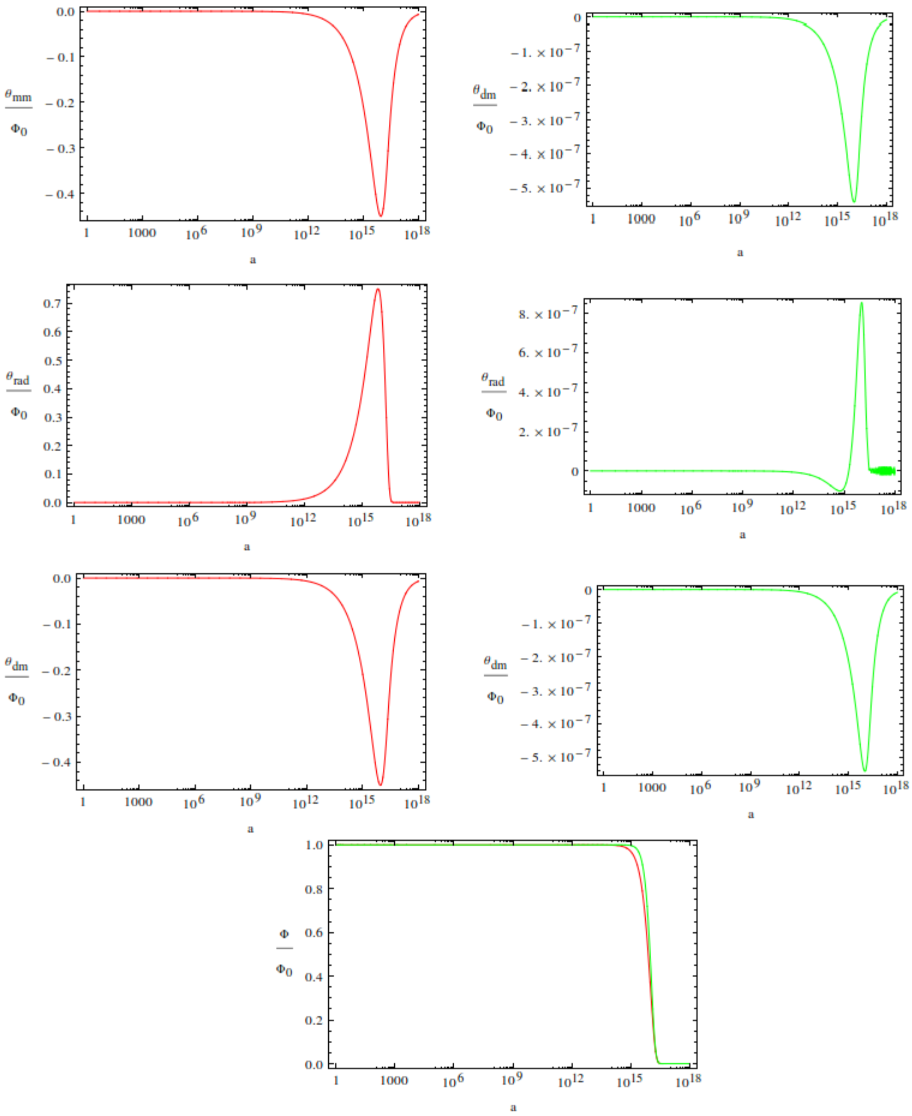

Fig. 7.7: Evolution in scale factor of the velocity components of the mother particle, radiation and daughter particle perturbations for two different scales. Red line corresponds to $k = 10^4 k_{RH}$ and green line to $k = 100 k_{RH}$. On the bottom it is pictured the evolution of the gravitational potential for the same pair of scales

Part IV

FINAL CONCLUSIONS

In this Thesis, non standard scenarios of Cosmology and particle physics have been proposed and studied. They were about currently hot topics in physics such as Inflation and the origin of the Baryon asymmetry and the observed astronomical objects. The results in this document can be viewed as one more contribution towards the knowledge of the physics in the very early Universe. Our aim has been then to explore and develop those topics that have not been fully considered in the literature regarding the physics before BBN, the widely accepted starting point of the Standard history of the Universe.

In the first chapter, the Inflationary potential triggered by a right handed neutrino condensate is considered. Such a condensate arises from a 4-Fermi interaction made up of right handed neutrinos following the same procedure given by Bardeen et *al* [6]. As a consequence, an effective Lagrangian for the condensate Φ is obtained, valid below a cut-off scale Λ. Moreover, the fact that right handed Neutrinos are Majorana particles that violate lepton number below a cut-off scale Λ' allows the phase θ of the condensate to develop a potential which tracks the "Natural" Inflationary potentials. It is important to highlight that such a form for the potential is still alive after Planck results [11, 57]. The inflationary behaviour of this dynamically generated potential depended on three parameters: v the scale related to the right handed neutrino mass, g and g', Yukawa couplings related to both Λ and Λ' respectively. In order to satisfy the experimental constraints on the amplitude of the primordial perturbations and spectral index, it was found the following bounds for our fundamental parameters

$$(g^3 g')^{1/2} \sim 10^{-5} \ , \tag{7.55}$$

$$v \sim M_{Pl} \ , \tag{7.56}$$

which tells us that our right handed neutrino mass is heavily pushed to the Planck scale.

In addition, the question of the generation of the baryon asymmetry was considered. Our theory consists only of the SM Lagrangian and the effective Lagrangian generated by the above mechanism. Therefore, the baryon asymmetry could only be produced in the scenario triggered by our condensate. We then showed that the appearance in the Lagrangian of an operator which couples the derivative of our inflaton to the lepton current yielded a net lepton number when the velocity develops a non zero expectation value. Such a mechanism can only work if lepton number and CPT are violated, where the latter one only occurs temporarily until the velocity approaches to zero. We then showed that the requirement of the baryon asymmetry allowed us to further constrain the parameters of our theory to

$$g' \simeq \quad 10^{-7} \tag{7.57}$$

$$g \simeq \quad 0.1 \tag{7.58}$$

which are reasonable given the fact that g' comes from an explicitly lepton symmetry breaking and one would expect it to be small in general.

In sum, the scenario studied in the first paper has the benefits of dynamically providing an inflationary framework and contributing with a mechanism for baryogenesis, with the only minimal ingredient of Right Handed Neutrinos.

The second and third paper have the same key element in common: the domination of non relativistic matter during the early times in the thermal history of the Universe. As it was stressed in the introductory chapters, the Standard Cosmology scenario applies to the history of the Universe from BBN onwards. It is usually assumed that a period of RD preceded the standard thermal history of the Universe, but this does not need to be true. Thermal Inflation [9] is a famous example of this. In both papers however, we followed the recipe for an early MD proposed in [8]. The realisation of such a scenario in this work comes from a thermal abundance of a heavy particle that at some point turned out to decouple from the thermal bath and dominate the energy density of the Universe. Additionally, one needs to allow a decay channel for this particle in order to link to the standard Universe by BBN. With these ingredients, the implications of such a scenario were studied on the electroweak phase transition and the growth of structures.

In paper II, we dealt with the nature of the electroweak phase transition under an early period of MD. We then derived the sphaleron bound, which constrains the survival of any baryon asymmetry at the EWPT, as a function of the Hubble rate H. We showed that the ratio between the Higgs vev $\langle\phi(T)\rangle$ and the critical temperature T_c could be relaxed if the Universe were dominated by matter at that time. We also showed that when plugged a Hubble rate in the radiation case, we recovered the usual sphaleron bound. Thus, we used this relaxation to study the baryon asymmetry generation at the EWPT in both SM and MSSM. We found that the SM is still unable to yield the baryon asymmetry of the Universe given the current bound on the Higgs mass. Furthermore, even if the strength of the phase transition happened to be first order enough, the CP violation in the SM is still too small to account for the asymmetry. On the other hand, the MSSM scenario is more interesting. We found that in the most favourable case, the sphaleron bound can be diminished to $\phi(T_c)/T_C \gtrsim 0.6$, which would be sufficient to open the parameter space window in the Light Stop Scenario, specially important after the last LHC results that are pushing the SUSY scale to large values. The main drawback that we found is that the component that dominated the Universe during the EWPT produced entropy with the consequent dilution of the baryon asymmetry. As a result of this, we argued that the baryogenesis mechanism in this scenario needs to be more efficient, a fact difficult not impossible to achieve. We finally showed in paper II a different realisation of the modification of the Hubble rate at the electroweak scale by resorting to braneworld Cosmologies. It was found that such a scenario can be easily achieved with reasonable matching temperatures between the braneworld and standard Scenario by BBN.

During the domination of Dark matter after the matter-radiation equality point, the lack of pressure allows the gravitational wells to begin to accrete matter and perturbations can grow. Similarly, we considered in paper III, the growth of structures during a MD epoch prior to BBN. Such a domination is conducted by a heavy non relativistic component, that we called 'mother', decaying likewise in matter and radiation in order for us to link with the Standard Big Bang scenario. We argued that the matter production, that was labelled as daughter, needs to play the role of Dark Matter if we want to maximise the effects on the growth of structures. Therefore, in order to get the right amount of dark matter today, we obtained an upper limit for the branching fraction (the amount of daughter matter from the decay of the mother) of $f_b \leq 10^{-6}$, which clearly disfavoured the creation of complex objects due to their evaporation from the entropy production. In order to avoid such a constraint, we proposed to add an annihilation channel into radiation for the daughter particle, which primarily affects the thermal history of the Universe shortly after the early MD epoch and allowed us to raise the upper bound of the branching fraction f_b, getting then a higher abundance of dark matter during that epoch. Under these conditions, we found that perturbations which entered the horizon early enough during this matter domination epoch can grow until becoming non-linear and start forming complex objects such as mini-halos, which in principle can be inferred through observational techniques [48, 49, 52–54]. We also estimated the size of these structures to belong to the miliparsec regime.

CONCLUSIONES FINALES

En esta Tesis, varios escenarios no estándar de Cosmología y física de partículas han sido propuestos para su estudio. Dichos escenarios tratan sobre temas considerados candentes tales como Inflación y el origen de la asimetría bariónica y de los objetos astronómicos que se pueden observar. Los resultados obtenidos en este trabajo pueden ser considerados como una pequeña contribución más hacia nuestro entendimiento de la física del Universo temprano. Nuestro propósito ha sido la exploración y desarrollo de ciertas cuestiones relacionadas con la física anterior a BBN y que no han sido completamente tratadas en la literatura.

En el primer capítulo, se estudió el potencial inflacionario originado a partir de un condensado de neutrinos dextrógiros. Dicho condensado surge siguiendo el procedimiento de Bardeen et *al* [6] a partir de una interacción de tipo 4-Fermi formada por neutrinos dextrógiros. Debido a esto, se obtiene un lagrangiano efectivo para el condensado Φ, que es valido para escalas de energía por debajo de una escala de corte Λ. Además, el hecho de que estos neutrinos dextrógiros sean de Majorana que violan el número leptónico por debajo de una escala de corte Λ' permite que la fase θ del condensado desarrolle un potencial que se asemeja a un potencial inflacionario de tipo "Natural". Es importante resaltar que dicho potencial sigue siendo válido después de los últimos resultados de PLANCK. El comportamiento inflacionario de este potencial generado dinámicamente dependía de tres parámetros: "v", la escala relacionada con la masa de los neutrinos, "g" y "g' ", que son operadores de Yukawa relacionados con las escalas Λ y Λ' respectivamente. Para satisfacer las restricciones experimentales en la amplitud de las perturbaciones primordiales y el índice espectral, se encontraron los siguientes límites sobre nuestros parámetros:

$$(g^3 g')^{1/2} \sim 10^{-5} \tag{7.59}$$

$$v \sim M_{Pl} \,, \tag{7.60}$$

que nos indican que nuestra masa de neutrino dextrógiro se encuadra fuertemente cerca de la escala de Planck.

Asimismo, también se consideró la cuestión de la generación de la asimetría bariónica. Nuestra teoría estaba compuesta sólo por el lagrangiano del Modelo Estándar y el lagrangiano efectivo que se generaba por el mecanismo anteriormente explicado. Es así entonces que la asimetría bariónica sólo podía ser generada en este escenario a través de nuestro condensado. Mostramos que la aparición en el lagrangiano de un operador que acoplaba la derivada de nuestro inflatón con la corriente leptónica producía un número leptónico neto cuando la velocidad desarrollaba un valor esperado en el vacío. Este mecanismo sólo puede funcionar si el número leptónico y CPT se violan, donde esta última violación se da de forma temporal hasta que la velocidad se aproxima a cero. Mostramos así que el requerimiento de que se produzca la correcta asimetría bariónica nos permitía restringir los parámetros de nuestra teoría a

$$g' \simeq 10^{-7} \tag{7.61}$$

$$g \simeq 0.1 , \tag{7.62}$$

valores que son razonables dado que g' proviene de una rotura explicita del numero leptónico y por tanto, uno esperaría que sea pequeña en general.

En conclusión, el escenario estudiado en el primer artículo tenía los beneficios de proporcionar dinámicamente un esquema inflacionario y contribuir con un mecanismo de bariogénesis, con el único ingrediente mínimo de Neutrinos dextrógiros.

El segundo y tercer artículo tienen el mismo elemento en común: la dominación de materia no relativista durante las épocas tempranas de la historia térmica del Universo. Como se enfatizó en los capítulos introductorios, el escenario estándar de Cosmología se aplica desde BBN en adelante. Se suele asumir que un periodo de radiación precedió a BBN en la historia térmica del Universo, pero esto no tiene por que ser cierto. Inflación termal [9] es un ejemplo famoso de esto. En ambos artículos, seguimos la receta dada en [8] para una dominación temprana de materia no relativista. La realización de dicho escenario en este trabajo venía de una abundancia térmica de una partícula muy masiva que en algún momento se desacopló del baño térmico y domino la densidad de energía del Universo. Adicionalmente, mostramos que uno necesitaba añadir un canal de desintegración para esta partícula para conectar este escenario con el estándar antes de BBN. Con estos ingredientes, se estudiaron las implicaciones de este escenario en la transición de fase electrodébil y el crecimiento de estructuras.

En el artículo II, tratamos la cuestión de la naturaleza de fase electrodébil durante un periodo dominado por materia no relativista. Calculamos así el límite de esfalerón, que restringe las posibilidades que cualquier asimetría bariónica tiene de sobrevivir durante la mencionada transición como función del rate de Hubble H. Mostramos que la relación entre el valor esperado en el vacío de Higgs $\langle\phi(T)\rangle$ y la temperatura critica T_c podía ser relajada si el Universo estuvo dominado por materia no relativista en ese instante. También mostramos que se recuperaba el límite de esfalerón usual cuando se usaba un Universo dominado por radiación. De esta forma, usamos esta relajación para estudiar la generación de la asimetría bariónica durante la transición de fase electrodébil dentro del Modelo Estándar y el MSSM. Encontramos que el Modelo Estándar seguía sin ser capaz de producir la asimetría bariónica dada la masa de Higgs. Además, incluso si la naturaleza de la transición de fase era suficientemente de primer orden, la violación de CP en el Modelo Estándar sería todavía insuficiente para producir dicha asimetría. Por otro lado, el escenario del MSSM es más interesante. Encontramos que en el caso más favorable, el límite de esfalerón puede disminuir a un valor $\frac{\phi(T_c)}{T_c} \geq 0.6$, que seria suficiente para extender el rango del espacio de parámetros permitidos en el escenario de Stops ligeros, escenario que cobra especial importancia tras los últimos resultados del LHC de que la escala de SUSY tiene que ser grande. El principal inconveniente que nos encontramos en nuestra propuesta fue que la producción de entropía por parte del componente que dominaba la transición de fase podía diluir la asimetría bariónica. Por ello, resaltamos que el mecanismo de bariogénesis necesita ser más eficiente en este escenario, un hecho difícil de alcanzar pero no imposible. Por último, también mostramos en el artículo II una diferente manera de modificar el rate de Hubble mediante las Cosmologías de mundo de branas. Se encontró entonces que este escenario puede ser fácilmente logrado con temperaturas razonables que unen el mundo de branas y el escenario estándar antes de BBN.

La falta de presión durante la dominación de Materia Oscura tras el punto de igualdad entre materia y radiación permite que los pozos de potencial empiecen a acretar materia y que las perturbaciones puedan crecer. De forma similar, consideramos en el articulo III el crecimiento de estructuras durante un periodo dominado por materia no relativista pero anterior a BBN. Esta dominación era llevada a cabo por una componente muy pesada, que nosotros llamamos “madre” y que decaía a su vez en materia y radiación con el fin de conectar con el escenario estándar del Big Bang. Argumentamos entonces que la producción de materia, a la que etiquetamos como materia “hija”, tenía que emplear el papel de Materia Oscura si queríamos recuperar los efectos del crecimiento de estructuras. Debido a esto, obtuvimos una cota superior en la fracción de desintegración $f_b \leq 10^{-6}$, que claramente desfavorecía la creación de objetos complejos dada su evaporación por la producción de entropía. Para solventar esto, propusimos añadir un canal de aniquilación de la partícula hija en radiación, que principalmente afectaba la historia térmica del Universo justo después de la época temprana de dominación de materia no relativista. Esto nos permitió subir la cota de f_b. Así, con estas condiciones encontramos que las perturbaciones que entraron durante esta época temprana de materia no relativista podían crecer hasta llegar al régimen no lineal y empezar a formar objetos complejos tales como los mini-halos, que en principio pueden ser inferidos y detectados a través de técnicas observaciones [48, 49, 52–54]. Por último, también estimamos que el tamaño de estas estructuras corresponderían al régimen de mili-parsec.

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