Deflationary Models Driven by Matter Creation

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Abstract

A nonsingular deflationary cosmology driven by adiabatic matter creation is proposed. In such a scenario there is no preinflationary stage as happens in conventional inflationary models. Deflation starts from a de Sitter spacetime characterized by an arbitrary time scale $H_I^{-1}$, which also pins down an initial value for the temperature of the universe. The model evolves continuously towards a slightly modified Friedman-Robertson-Walker universe. The horizon and other well known problems of the standard model are then solved but, unlike in microscopic models of inflation, there is no supercooling and subsequent reheating. Entropy generation is concomitant with deflation and if $H_I^{-1}$ is of the order of the Planck time, the present day value of the radiation temperature is deduced. It is also shown that the “age problem” does not exist here. In particular, the theoretically favored FRW flat model is old enough to agree with the observations even given the high values of $H_o$ suggested by recent measurements.

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1 Introduction

The problem of cosmological particle creation and entropy generation is presently a very active field of research. It is now widely believed that matter and radiation need to be created, at least in the very early universe, in order to overcome some difficulties presented by hot big-bang cosmology [1-17].

From a microscopic point of view, such processes were first investigated by Parker and collaborators [1] by considering the Bogoliubov mode-mixing technique in the context of quantum field theory in curved space-time [2]. Despite being rigorous and well-motivated, those models were never fully realized probably, due to the lack of a well-defined prescription of how matter creation is to be incorporated in the classical Einstein field equations (EFE). This issue has been further explored in all variants of the inflationary scenario [3, 4]. Usually, the amplification of zero-point fluctuations of quantum fields by the interaction with the background curvature and the corresponding matter creation rates are studied in close analogy with some processes in quantum optics [4].

The consequences of matter creation have also been macroscopically investigated mainly as a byproduct of bulk viscosity processes near the Planck era as well as in the slow-rollover phase of the new inflationary scenario [5-9]. More recently, the first self-consistent macroscopic formulation of the matter creation process was put forward by Prigogine and coworkers [10] and somewhat clarified by Calvão et al. [11, 12]. It was argued that matter creation, at the expenses of the gravitational field, can effectively be discussed in the realm of nonequilibrium thermodynamics. In comparison with the standard equilibrium equations, the process is described by two new ingredients: a balance equation for the particle number density and a negative pressure term in the stress tensor. Such quantities are related to each other in a very definite way by the second law of thermodynamics. In particular, the creation pressure depends on the creation rate and may operate, at level of the EFE, to prevent either a spacetime singularity [10, 16] or to generate an
early inflationary phase \[13, 18\].

As it appears, such an approach may be crucial in establishing a pattern for the thermodynamics of a viable quantum-mechanical matter creation process since it seems to incorporate, in a rather simple scheme, the backreaction contributions. In this formulation, the laws of nonequilibrium thermodynamics were used since the very beginning, thereby leading to definite relations among classical quantities and, more important, establishing the temperature evolution law for the created particles. In this way, the fragile semiclassical analysis needed by the former models of matter creation as well as by some analog treatments, such as the one represented by bulk viscosity mechanisms, lose significance \[13, 14\]. The stress tensor for matter creation now has a well-defined expression which depends on the usual physically meaningful quantities.

Inflation, on the other hand, has also been invoked as a key aspect of any viable cosmological scenario, since in its latest versions it could explain several puzzles of the visible universe: the horizon problem, the entropy of the cosmic microwave background (CMBR) and the causal generation of seeds to form the large-scale structure\[15\]. It is interesting to note, however, that the inflationary scenario is not free of their own drawbacks, which seem to be a strong evidence that such a mechanism is probably only a piece of a more complete cosmological theory. For instance, in some inflationary variants, there is a preinflationary stage in which the universe emerges from a radiation-dominated FRW phase and, due to the vacuum dominance, enters in the de Sitter epoch at a given critical temperature. This means that inflation does not evade the singularity problem, where matter and radiation were \textit{ab initio} created in a single event at \(t = 0\). In fact, as shown by Borde and Vilenkin\[21\], if some reasonable physical conditions are satisfied such models must necessarily possess an initial singularity. On the other hand, due to adiabatic expansion of the de Sitter phase, the universe undergoes a supercooling in which the temperature decreases by a factor of \(10^{-28}\) in order to solve the horizon problem. Subsequently, it is reheated through a
highly nonadiabatic process during the coherent oscillations of the inflaton field\[3\]. In this connection, it is worth mentioning that there is neither a coherent kinetic theoretical treatment nor a nonequilibrium thermodynamics formulation describing phenomena very far from equilibrium. Another well-known difficult of inflation is related to the “age problem”: the theoretically favored FRW flat model has an expansion age of \(\frac{2}{3} H^{-1}_o\) which by the latest measurements \[23, 24\] gives only 8.3 Gyr or less.

In this context, the present article discusses a class of deflationary cosmologies endowed with “adiabatic” matter creation. By deflation we mean inflation without a preinflationary stage i.e. the beginning of the universe in the remote past is like a pure de Sitter spacetime, irrespective of the curvature parameter. The “adiabatic” condition characterize thermodynamically the creation process. In this way, particles (and consequently entropy) are continuously generated, however, the specific entropy per particle remains constant during the process\[11, 12\]. In the case of photons this also means that the equilibrium relations are preserved\(n \sim T^3, \rho \sim T^4\) and that the photon spectrum is compatible with the present isotropy of the cosmic background radiation\[19, 20\].

As we shall see, due to the existence of a long period of deflation, the aforementioned problems are solved in a natural way. Unlike inflation, this model predicts that the de Sitter phase slowly cools down, smoothly connecting to a FRW-like evolution scheme. One clear advantage of such a model is that the universe is always in quasi-equilibrium, as opposed to having several “hiccups” of extremely out-of-equilibrium phase transitions. In particular, the rapid supercooling followed by the dramatic reheating existing in the microscopic models have been avoided. The model is nonsingular with the cosmic history starting from a de Sitter state characterized by an arbitrary time scale \(H^{-1}_I\). If \(H^{-1}_I\) is of the order of the Planck time, the present day temperature of the CMBR can be derived. It is also shown that the age of the universe is large enough to agree with observations even given the high value of \(H_o\) suggested by the latest measurements.
2 Basic Equations

We start with the homogeneous and isotropic FRW line element
\[ ds^2 = dt^2 - R^2(t) \left( \frac{dr^2}{1 - kr^2} + r^2 d\theta^2 + r^2 \sin^2(\theta) d\phi^2 \right), \]
where \( R \) is the scale factor and \( k = 0, \pm 1 \) is the curvature parameter. Throughout we use units such that \( c = 1 \).

In that background, the nontrivial EFE for a fluid endowed with matter creation and the balance equation for the particle number density can be written as \cite{12, 13, 18}

\[ 8\pi G \rho = 3 \frac{\dot{R}^2}{R^2} + 3 \frac{k}{R^2}, \]
\[ 8\pi G (p + p_c) = -2 \frac{\dot{R}}{R} - \frac{\dot{R}^2}{R^2} - \frac{k}{R^2}, \]
\[ \frac{\dot{n}}{n} + 3 \frac{\dot{R}}{R} = \frac{\psi}{n}, \]
where an overdot means time derivative and \( \rho, p, p_c, n \) and \( \psi \) are the energy density, thermostatic pressure, creation pressure, particle number density and matter creation rate, respectively. The creation pressure \( p_c \) is defined in terms of the creation rate and other thermodynamic variables. In the case of adiabatic matter creation, it is given by
\[ p_c = -\frac{\rho + p}{3nH}\psi, \]
where \( H = \dot{R}/R \) is the Hubble parameter.

By combining the usual “gamma-law” equation of state
\[ p = (\gamma - 1) \rho, \]
with the EFE, it is easily seen that the evolution equation for the scale function can be cast in the form below\cite{18}
\[ R\ddot{R} + \left( \frac{3\gamma - 2}{2} - \frac{\gamma \psi}{2nH} \right) \dot{R}^2 + \left( \frac{3\gamma - 2}{2} - \frac{\gamma \psi}{2nH} \right)k = 0 \quad . \]  

(7)

Consider now the following matter creation rate:

\[ \psi = 3\beta nH + 3(1 - \beta)n \frac{H^2}{H_I} \quad . \]  

(8)

The first term in the right-hand-side (RHS) of above equation is the same creation rate considered in Ref.\[18]\) (from now on referred to as paper I). The second one is a correction term of the order of \(\frac{H}{H_I}\). For \(\beta = 0\), Eq.(8) reduces to the matter creation rate considered by Lima and Germano\[13\], the consequences of which have been analysed only in the flat case. The main difficulty of such a scenario is closely related with the predicted small time interval elapsed during the FRW phase. As shown in \[18\), the presence of the \(\beta\) parameter is crucial to solve the “age problem” since it parametrizes the extent to which the model departs from FRW phase at late stages. The arbitrary time scale \(H_I^{-1}\) characterizes the initial de Sitter phase and, together with the \(\beta\) parameter, will presumably be given by some fundamental model of matter creation. At late times e.g., for \(H \ll H_I\), the first term on the RHS of (8) is dominant. Hence, the model discussed here may also be viewed as an early phase of the “adiabatic” matter creation scenario proposed in the paper I. Using the above expression we can recast Eq.(7) in the FRW-type form:

\[ R\ddot{R} + \left( \frac{3\gamma_\ast - 2}{2} \right) \dot{R}^2 + \left( \frac{3\gamma_\ast - 2}{2} \right)k = 0 \quad , \]  

(9)

where \(\gamma_\ast\) is an effective time-dependent “adiabatic index” given by

\[ \gamma_\ast = \gamma(1 - \beta)(1 - \frac{H}{H_I}) \quad . \]  

(10)

For either \(H = H_I\) or \(\beta = 1\), equation (10) gives \(\gamma_\ast = 0\), and (8) reduces to

\[ R\ddot{R} - \dot{R}^2 - k = 0 \quad , \]  

(11)
which yields the well known de Sitter solutions

\[
R(t) = \begin{cases} 
H_I^{-1} \cosh(H_I t) & k = +1, \\
R_0 e^{H_I t} & k = 0, \\
H_I^{-1} \sinh(H_I t) & k = -1.
\end{cases}
\] (12)

Hence, unlike in the standard FRW model, the present scenario begins in a pure nonsingular de Sitter vacuum with Hubble parameter \( H = H_I \) regardless the value of the curvature parameter. Accordingly, equation (2) gives \( \rho = \frac{3H_I^2}{8\pi G} \) as it should be due to the symmetries of the de Sitter spacetime. Analytically, the ansatz (8) can be viewed as the simplest matter creation law which destabilizes the initial de Sitter configurations given by (12). In this way, the initial evolution is such that the singularity, flatness and horizon problems may be simultaneously eliminated. All these asymptotic solutions have constant curvature, and are unstable in the future. Of course, closed (\( k = 1 \)) solutions are not of the “bouncing” type, rather the universe begins its evolution from a closed de Sitter universe.

In the opposite limit, \( H << H_I \), Eq.(10) reduces to \( \gamma_* = \gamma(1 - \beta) \) with (9) taking the following form

\[
R \ddot{R} + \Delta \dot{R}^2 + \Delta k = 0,
\] (13)

the first integral of which is

\[
\dot{R}^2 = \frac{B}{R^2\Delta} - k,
\] (14)

where \( \Delta = \frac{3\gamma(1-\beta)^2-2}{2} \) and \( B \) is a positive constant (see eq.(2)).

Using (14) one may express the energy density as well as the pressures (\( p \) and \( p_c \)) as functions solely of the scale factor \( R \) and of the \( \beta \) parameter. In fact, inserting (14) into (2), one obtains

\[
\rho = \rho_o \left( \frac{R_0}{\dot{R}} \right)^{3\gamma(1-\beta)},
\] (15)

where \( \rho_o = \frac{3B}{8\pi G R_o^{3\gamma(1-\beta)}} \). The above equation shows that the densities of radiation and dust scale, respectively, as \( \rho_r \sim R^{-4(1-\beta)} \) and \( \rho_d \sim R^{-3(1-\beta)} \).
Hence, the transition from a radiation to a dust dominated phase, in the course of the expansion, happens exactly as in the standard model. Of course, a more simplified scenario may also be implemented. For instance, if one assumes that only photons are produced such a transition is even more smooth than in the FRW case, since the $\beta$ factor for the matter energy density ($\gamma = 1$) in (15) must be suppressed. Another possibility is to consider that the $\beta$ parameter is always nonnull, however, assuming different values when the universe evolves from radiation to a matter dominated phase.

Summarizing, the cosmic history in the deflationary scenario proposed here proceeds in three stages. Firstly, there is a natural transition from an early de Sitter phase to a slightly modified FRW stage irrespective of the values of $k$ and $\gamma$. For $\gamma = \frac{4}{3}$, the model enters smoothly in the radiation dominated phase providing an interesting solution to the “exit problem” of the old inflationary scenario. Subsequently, the transition to the present dust phase is completed in the same fashion of the standard model. Indeed, the quasi-FRW stage, which is defined by condition $H \ll H_I$ correspond, at the level of the matter creation rate (8), to considering only the first term on the RHS of Eq.(8). This case has been studied in detail in paper I. As shown there, the $\beta$ parameter plays an important role during that phase since it increase the age of these models. Note also that all curvature invariants are bounded for any instant of time, and as a consequence the geodesics of comoving observers are unbounded. It thus follows that such spacetimes evolve from a nonsingular initial state.

For simplicity and also to obtain an exact description we now consider $k = 0$ as preferred by inflation. In terms of the Hubble parameter we can rewrite (9) as

$$\dot{H} + \frac{3\gamma(1-\beta)}{2}H^2(1 - \frac{H}{H_I}) = 0 \quad ,$$

which is just the equation of motion for the flat vacuum decaying model of Ref.\[22\]. The solution of (16) can be written as
\[ H = \frac{H_I}{1 + CR^{3\gamma(1-\beta)/2}} , \tag{17} \]

where \( C \) is a \( \gamma \)-dependent integration constant. Note that \( H = H_I \) is a special solution of (16) describing the de Sitter spacetime for any value of \( \gamma \). Such a solution is clearly unstable with respect to the critical value \( C = 0 \). For \( C > 0 \), the universe starts without singularity and evolves continuously towards a quasi-FRW stage, \( R \sim t^{2/3\gamma(1-\beta)} \), for large values of the cosmological time. It is also very easy to show that such a universe is horizon free. in fact, a light pulse beginning at \( t = -\infty \) will have traveled by the time \( t \) a physical distance \( l_H(t) = R(t) \int_{-\infty}^{t} \frac{dt'}{R(t')} \). From (17) we may write

\[ l_H(t) = H_I^{-1} R(t) \int_{0}^{R(t)} \frac{(1 + CR^{3\gamma(1-\beta)/2})dR}{R^2} . \tag{18} \]

Since the above integral diverge at the lower limit, the model is free of horizons e.g., light signals could have traveled to infinite distance at any \( t \), since the universe came into existence at coordinate time \( t = -\infty \) thereby, allowing the interactions to homogenize the whole universe.

The dynamic qualitative behavior sketched above may be described exactly. To show this, it proves convenient to compute \( C \) in terms of the present day parameters \( H_o, R_o \) and also of \( H_I \). From (17) one reads off

\[ C = \left( \frac{H_I - H_o}{H_o} \right) \frac{1}{R_o^{3\gamma(1-\beta)/2}} \]

and by recasting this constant in terms of

\[ R_\ast = R_o \left( \frac{H_o}{H_I - H_o} \right)^{2/3\gamma(1-\beta)} \]

which is determined by taking \( H(R_\ast) = H_I/2 \), one can write the integral of (17) in the convenient form

\[ H_I t = \ln \left( \frac{R}{R_\ast} \right) + \frac{2(H_I - H_o)}{3\gamma(1-\beta)H_o} \left( \frac{R}{R_0} \right)^{3\gamma(1-\beta)/2} \]. \tag{20}
Once $H_I$ is given, scale $R_\gamma(\gamma, \beta)$ automatically defines, for each model, the end of the deflation and the beginning of the quasi-FRW phase, as seen in Fig. 1. At early times ($R \ll R_\gamma$), when the logarithmic term is dominant, we obtain

$$R \simeq R_\gamma e^{H_I t},$$

in accordance with Eq. (12). At late times ($R \gg R_\gamma, H \ll H_I$) the expression reduces to the second term in (20)

$$R \simeq R_0 \left[ 3\gamma(1-\beta) \frac{H_o t}{2} \right]^{2/3\gamma(1-\beta)}.$$  (21)

The above expression shows us how the $\beta$ parameter may be useful in reconciling the theoretically favored $\Omega_o = 1$ with the latest measurements of the Hubble parameter [23, 24]. In a matter dominated universe ($\gamma = 1$), the time interval $\Delta t = t_o$ elapsed in the quasi-FRW stage is $t_o \sim \frac{2H_o^{-1}}{3(1-\beta)}$, when in the standard flat model ($\beta = 0$), one would obtain exactly $\frac{2}{3}H_o^{-1}$ (see also Fig. 1). As remarked earlier, the arbitrary time scale $H^{-1}_I$ furnishes the greatest value of energy density ($\rho = \frac{3H^2}{8\pi G}$). In addition, it is readily seen that the maximum matter creation rate, $\psi_I = 3n_I H_I$, also occurs in the de Sitter phase. In fact, by inserting (8) into (4), the balance equation for the particle number density assumes the form

$$\dot{n} + 3n(1-\beta)H(1 - \frac{H}{H_I}) = 0,$$  (22)

and from (17), the solution of the above equation may be cast as

$$n = n_I [1 + \left(\frac{R}{R_\gamma}\right)^{3\gamma(1-\beta)/2}]^{-\frac{2}{\gamma}}.$$  (23)

In the de Sitter phase, $R \ll R_\gamma$, the above equation yields a constant particle number density, $n \sim n_I$, and as a consequence the net number of particles $N$ grows proportional to $R^3$. In the opposite regime, $R \gg R_\gamma$ ($H \ll H_I$), $n$ scales with $R^{-3(1-\beta)}$ as expected for the quasi-FRW stage (see paper I). In what follows, the above results will be extensively used.
3 Thermodynamic Behavior

Let us now discuss some thermodynamic features of the deflationary scenario proposed in the later section. For adiabatic matter creation the temperature law and the rate of variation of the specific entropy are given by\([12, 13]\)

\[
\frac{\dot{T}}{T} = \left( \frac{\partial p}{\partial \rho} \right) \frac{\dot{n}}{n},
\]

(24)

\[
\dot{\sigma} = 0,
\]

(25)

where \(T\) is the temperature and \(\sigma\) is the specific entropy (per particle). Using the \(\gamma\)–law equation of state, a straightforward integration of (24) yields

\[
T = T_I \left( \frac{n}{n_I} \right)^{\gamma^{-1}},
\]

(26)

or still, inserting (23)

\[
T = T_I \left[ 1 + \left( \frac{R}{R_s} \right)^{3\gamma(1-\beta)/2} \right]^{-2\frac{\gamma-1}{\gamma}},
\]

(27)

where \(T_I\) is the (maximal) temperature at de Sitter phase. Due to irreversible matter creation, the expansion proceeds isothermically during the de Sitter phase \((R < R_s)\), thereby avoiding the extremely rapid supercooling and the subsequent dramatic reheating that must take place at \(R \simeq R_s\) in all inflationary scenarios\([3]\). The temperature decreases continuously in the course of the expansion. At late stages, when the universe enters in the quasi-FRW phase, the temperature scales as \(T \sim R^{-3(1-\beta)}\), as expected (see paper I). For \(\beta = 0\), the result of Ref.\([13]\) is recovered, with the universe evolving towards the standard FRW model. Now, recalling that \(R_s\) was determined in terms of \(H_I, H_0, R_0\) and \(\beta\) by (19), the above temperature evolution law depends only on three unknown quantities, namely: Two initial conditions \(T_I\) and \(H_I\), plus the \(\beta\) parameter. Note, however, that the \(\beta\)-dependence in (27) cancels at \(R = R_o\) thereby, reducing our problem to the right choice of the initial
conditions. Since the model starts as a de Sitter spacetime, the most natural choice is to define $T_I$ as the Hawking temperature

$$T_I = \frac{\hbar H_I}{2\pi k_B}, \quad (28)$$

Of course, such a choice is dictated only by the overall existence of an initial de Sitter phase so that it could be applied irrespective of the value of the curvature parameter. The arbitrary time scale of the de Sitter state, $H_I^{-1}$, has not been fixed by the model. For consistency, since we are working with a classical description, $H_I^{-1}$ must be chosen in such way that (28) becomes of the same order or smaller than Planck temperature. Indeed, in the framework of quantum cosmology, many authors have suggested that the spontaneous birth of the universe leads naturally to a de Sitter stage with $H^{-1} \sim t_{pl}$ or equivalently an energy density $\rho \sim \rho_{pl}$ (see for example [25]). It is remarkable that such a choice, say, $H_I = \frac{2t_{pl}}{\alpha}$, where $\alpha$ is a pure number, allow us to estimate the present value of the temperature of CMBR. In fact, from (28), the initial temperature of the universe is close to Planck temperature

$$T_I = \frac{1}{\alpha k_B} \sqrt{\frac{\hbar}{G}} \sim \frac{1.4 \times 10^{32}}{\alpha} K. \quad (29)$$

Now, taking $\gamma = 4/3$ and substituting $R_*$ into (27) we may write

$$T = T_I[1 + \frac{H_I - H_o}{H_o}(R/R_o)^{2(1-\beta)}]^{-\frac{1}{2}}, \quad (30)$$

and since $H_I >> H_o$, replacing (29) and using the previously defined value of $H_I$, we obtain at the present time

$$T_o = \frac{1.4 \times 10^{32}K}{\sqrt{\alpha}} \left(\frac{H_o t_{pl}}{2\pi}\right)^{1/2}, \quad (31)$$

where $T_o$ is the temperature at $R = R_o$. The above expression may be used either to estimate $T_o$ if the value of $H_o$ is given or to compute $H_o$ using the well established measurements of $T_o$. For example, assuming that $H_o$ is centered at $80 \text{ km s}^{-1} \text{ Mpc}^{-1}$ as claimed by Freedman et al. [24], this means
that $H_o \approx 2.7 \times 10^{-18} s^{-1}$ and since $t_P \sim 5.4 \times 10^{-44}s$, we obtain $T_o \sim \frac{21K}{\sqrt{\alpha}}$. Therefore, taking $\alpha$ of the order of 50, it follows that $T_o \sim 2.8K$ and from (28), the initial temperature of the universe is 50 times smaller than Planck temperature.

Let us now discuss the entropy production. By definition, the specific entropy (per particle) is $\sigma = \frac{S}{N}$, it thus follows from (24) that

$$\frac{\dot{S}}{S} = \frac{\dot{N}}{N},$$

and since $\dot{N} \geq 0$ in all stages, the total entropy increases in the course of the evolution. Unlike of all microscopic variants of inflation, entropy generation in our scenario is concomitant with the overall deflationary process. To compute the net entropy in terms of the scale function, it proves convenient to consider the entropy density $s_I$ of the de Sitter phase as an initial condition. In addition, instead of integrating (32), it is much simpler to consider (23) together with the constant value of $\sigma = \sigma_I$, to obtain

$$S = \frac{s_I R^3}{\left[1 + \left(\frac{R}{R^*}\right)^{3\gamma(1-\beta)/2}\right]^{\frac{4}{7}}}.\quad (33)$$

Hence, during the de Sitter phase ($R << R_s$) entropy scales as $S \sim R^3$, whereas at late stages $S \sim R^{3\beta}$ as it should (see paper I). In Fig. 2 we have plotted energy density $\rho$, total entropy $S$ and temperature $T$ in the early stages of a radiation-dominated model with $\beta = 0.25$.

For $\beta = 0$, (33) reduces to equation (27) of Ref.[13]. In this case, the entropy is constant at late times as expected for an adiabatic FRW type expansion. Now, inserting the value of $R_s$ given by (19) with $H_I >> H_o$, taking $\gamma = 4/3$ and $R_o \sim H^{-1}_o$, to a high degree of accuracy (33) yields $S_o \sim s_I \left(\frac{H_o}{H_I}\right)^{\frac{3}{2}}$. Of course, the value of $s_I$ (or $n_I$) is uniquely defined in terms of $T_I$. Indeed, since the equilibrium relations are maintained, $\rho \sim T^4$ and $n \sim T^3$ (see (26) and paper I), it follows that $s_I = \frac{4}{3} a T^3_I$, where $a = \frac{\pi^2}{15}$ is the radiation equilibrium constant ($\hbar = k_B = 1$). Finally, by replacing the value
of $H_I$, the present radiation entropy in our horizon may be expressed as

$$S_o \approx \frac{3.10^{-2}}{(\alpha H_o t_P)^{\frac{3}{2}}}.$$  \hspace{1cm} (34)

Therefore, as we did in the case of temperature, taking $\alpha = 50$ and $H_o = 80 \text{ km s}^{-1} \text{ Mpc}^{-1}$, we obtain for the present dimensionless radiation entropy $S_o \sim 2.9 \times 10^{87}$.

4 Conclusion

A class of deflationary cosmologies driven by “adiabatic” matter creation has been proposed. In these models, the beginning of the universe is a pure de Sitter state supported by the matter creation process. Subsequently, the universe evolves smoothly towards a radiation dominated phase as happens in the standard model. In this way, both the singularity and horizon problems have been naturally solved. As a matter of fact, deflationary spacetimes seem to be a rather generic solution of the EFE since it can also be generated by quite different mechanisms like bulk viscosity[9] and vacuum decaying energy density[22].

Unlike all variants of inflation, in such models there is no supercooling or reheating. The temperature of radiation diminishes slowly from its greatest value during the initial de Sitter configuration until the present 2.8K. Simultaneously, the entropy increases continuously to reach $S_o \sim 10^{87}$ at the present phase. In addition, due to persistent matter creation the age of the universe also fits well with a higher value of $H_o$. More important, such results were obtained using one and the same initial condition. In this connection, we remark that the initial Hawking-Gibbons temperature given by (28) is not related here with the horizon temperature of a vacuum de Sitter spacetime. Rather, this de Sitter phase is supported by the negative creation pressure and for $\gamma = \frac{4}{3}$ space is filled with radiation. At late times, the model evolves to the scenario proposed in Ref.[18], thereby solving the “age problem”. As
argued there (see also [19, 20]), such a scenario is also compatible with the constraints from the cosmic background radiation, and a crucial test for these models is provided by the measurements of the temperature of the universe at high redshifts.

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Figure Captions

Figure 1:

**Fig. 1** - Scale factor as a function of time for $\gamma = 4/3$ (radiation) and $\beta = 0$ (left curve), $\beta = 0.25$, $\beta = 0.5$ and $\beta = 0.75$ (right curve). For $\beta > 0$, the universe deflates smoothly from an initial de Sitter phase ($R \ll R_*$) to a quasi-FRW stage ($R \gg R_*$). The dashed curve depicts scale factor of a FRW universe starting at $t = 0$. When $\beta = 0$ we have no matter creation at the present time, and for $\beta \to 1$ the model inflates at all times powered by a steady creation rate. The physically meaningful ages are easily seen to be the time elapsed from $t = 0$ to the moment when a curve crosses a fixed scale ($R = R_*$). As expected, models with larger matter creation at late times (larger $\beta$) are older.

Figure 2:

**Fig. 2** - Energy density, total entropy and temperature of radiation as a function of the scale factor (normalized to $R_*$) in a scenario with $\beta = 0.25$. Notice the nearly constant energy density and temperature, but a huge growing entropy during the deflationary process. As soon as the scale $R = R_*$ is reached, regimes change and the FRW-like epoch, characterized by falling energy density and temperature, starts. The model does not evolve adiabatically. The total entropy always increases, although moderately nowadays.