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BULK-VISCOSITY-DRIVEN ASYMMETRIC INFLATIONARY UNIVERSE

by

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ABSTRACT

A primordial net bosonic charge is introduced in the context of the bulk-viscosity-driven inflationary models. The analysis is carried through a macroscopic point of view in the framework of the causal thermodynamic theory. The conditions for having exponential and generalized inflation are obtained. A phenomenological expression for the bulk viscosity coefficient is also derived.

Key-words: Early universe; Inflation; Bulk viscosity; Phase transition.

I INTRODUCTION

Recently, an increasing attention has been paid on the bulk viscosity phenomenon associated with the inflationary models of the universe[1-6]. In fact, the effect of bulk viscosity in an expanding universe is to reduce the equilibrium pressure. So, it is natural to ask if this effect could be strong enough to make the effective pressure negative. As is well known, this is the key condition for inflation.

Thermodynamic states with negative pressure are metastable and are not excluded by any law of nature. In general, these states are connected with phase transitions (for example in an overheated Van der Waals liquid[7]) and for certain physical systems the occurrence of negative pressure seems to be inevitable[8,9,10]. These systems are hydrodynamically unstable for bubbles and cavities formation and spontaneous collapse could also be expected[7]. However, in the cosmological context, some new features must be added. In fact, as was shown by Whittaker[11], in a stressed selfgravitating fluid described by general relativity, the pressure also contributes to the effective gravitational mass. If this contribution is negative it will act repulsively accelerating the cosmic expansion. Therefore, a fluid out of thermodynamic equilibrium with negative effective pressure provides an alternative mechanism for inflation.

More recently some authors[12] claimed that the bulk viscosity could not drive inflation. They remarked that in the framework of kinetic theory, the pressure is always non negative. Hence, it is important to know under which conditions the kinetic approach can be applied to the early universe, in particular, at the epoch of the spontaneous symmetry breaking (SSB) of the grand unified theory (GUT)[13].

A coherent kinetic-theoretic treatment can be accomplished only if the system in question present a sufficient dilution degree. The validity of such approximation depends on the ratio λ/L, where λ is the mean free path and L is the mean interparticle distance. By using that $\lambda \approx (\sigma n)^{-1}$ and $L \approx n^{(-1/3)}$, where σ is the interaction cross-section and n is the number density, this ratio is easily estimated at the GUT era. In virtue of the asymptotic freedom we have $\sigma \approx \alpha^2/T^2$, where $\alpha \approx 1/40$ is the unified coupling strength. Now, by reasons that will be presented bellow we assume the existence of a primordial Bose-Einstein condensate. So, $n \approx q + N_* T^3$ where q is the net bosonic charge density and $N_* = (\pi^2/30)(\Sigma g_b + 7/8 \Sigma g_f)$ where g_b and g_f denote the number of effectively massless bosonic and fermionic degrees of freedom, respectively. The term "charge" is simply the difference between bosons and antibosons number[14,15]. Therefore, it follows that

$$\frac{\lambda}{L} \approx \alpha^{-2} N_*^{-2/3} \left(1 + \frac{q}{N_* T^3}\right)^{-2/3} \qquad (1.1)$$

Now considering that at $T \approx 10^{14}$ GeV the total number of particle species is ≈ 150, we have $N_* \approx 50$. Then, if the universe is symmetric ($q \simeq 0$), $\lambda/L \approx 100$ and the dilute gas approximation is a good one, but for instance, if $q \approx 10^6 T^3$ then $\lambda/L \approx 10^{-1}$, the continuum GUT is dense and, in consequence, the kinetic approach cannot be applied.

Haber and Weldon[15,16] showed that the condensation critical temperature in a relativistic ideal Bose-Einstein gas is given by

$$T_c = (3q/m)^{1/2} , \qquad (1.2)$$

where m is the mass of the particles. In fact, the formula above is valid only for $T \gg m$ but it will be sufficient for our qualitative arguments. These authors also observed that if the bosons are massless the critical temperature T_c is infinite and so all the net charge will be in the Bose-Einstein condensed ground state. This result is important in what will be proposed ahead. Note that before the GUT phase transition the leptoquark gauge bosons are massless, thus all bosonic charge excess will be in the condensed phase. So, even when the charge density is high this avoids the possibility of a superclosed universe. During the phase transition these gauge bosons acquire a mass of order of the GUT energy scale. In consequence, the critical temperature T_c in equation (1.2) will fall to some finite value depending on the magnitude of the primordial charge. For a high value of this charge, the critical temperature T_c can be higher than the universe temperature T. So, just after the onset of the phase transition it is possible for the superheavy bosons to be in the condensed phase. In this case, another kind of phase transition should occur in virtue of the cosmic expansion. Note from (1.2) that T_c scales with $R^{-3/2}$ whereas in an adiabatic expansion, the universe temperature T scales with R^{-1} (R is the scale factor). Therefore, as the expansion proceeds, T_c decreases faster than T and the decondensation process will be in course. This second phase transition could occur at the end of inflation and the possible formation of vortex lines (as in superfluid helium) could be relevant for the structure of the universe[17].

If the charge excess is associated with the Higgs particles some aspects of the above discussed qualitative picture must be modified. In this case, for example, Bose-Einstein condensation and spontaneous symmetry breaking can occur simultaneously but independently. Moreover, studying Bose-Einstein condensation in the Weinberg-Salam model Kapusta[18] showed that the transition temperature is raised. In principle, analogous results could be derived for others gauge theories. This would be interesting for our scenario since the irreversibilities are expected to be more relevant if the GUT scale parameter is higher[1].

Our present knowledge of the GUT continuum is still quite limited, thus a thermodynamic approach might give a major flexibility to the model. Details of the microscopic theory will not be taken into consideration in this paper.

In a previous paper[4] some results uniting inflation and bulk viscosity were obtained in the framework of a "quasi-stationary" or first order relativistic theory of dissipative processes. This terminology has been largely adopted because the quadrivector entropy flux contains only first order terms in deviations from equilibrium. This paper made use of the Eckart-Weinberg's[19,20] formulation (Landau and Lifshitz[21] approach is also included in the above category). As is well known, these theories contain several undesirable features. They lead to parabolic differential equations and so admit superluminal velocities for heat flow and viscosity propagations. In addition, the theory is unstable and there is no well posed initial value problem for rotating fluids[22]. Now, besides introducing a chemical potential associated with a possible bosonic primeval charge, we reanalyze the bulk-viscosity-driven inflationary scenario in the context of a non-stationary (or transient) relativistic thermodynamic theory. This theory was developed by Müller[23] and Israel[24] and solves the above cited problems present in the first order theory. More details will be given in section II.

We outlined this paper as follows: Section II establishes the conditions for having exponential or generalized[25,26] inflation and the framework of our approximations are presented. In section III we obtain from a thermodynamical point of view the expression for the bulk viscosity coefficient, the entropy production and the duration of inflation.

II. CONDITIONS FOR INFLATION

We start supposing that ab initio the universe had a net bosonic charge. Moreover, as is usual in the inflationary models, we assume that: First, some region of the universe with the size about the horizon distance was hot ($T \geq T_{GUT}$) and cooled to the GUT critical temperature before recollapsing. Second, this region was sufficiently homogeneous and isotropic such that the Robertson-Walker (RW) metric is a good approximation.

In the early times the contributions of the spatial curvature were negligible then, to describe the geometry of the model during the inflation, we can use the flat RW metric

$$ds^2 = -dt^2 + R(t)^2\,(dx^2 + dy^2 + dz^2), \tag{2.1}$$

where R is the scale factor.

We will work in the framework of a charged relativistic simple fluid[27]. The functional dependence of the entropy flux vector S^α is

$$S^\alpha = S^\alpha(T^{\mu\lambda}, J^\mu) \tag{2.2}$$

where $T^{\mu\lambda}$ is the energy momentum tensor and J^μ is the charge flow vector defined by

$$J^\mu = qu^\mu \tag{2.3}$$

where u^{μ} is the velocity of the comoving observer. The Lorentz invariance of the theory places a severe restriction on the form of $T^{\mu\lambda}$ in local equilibrium, that is:

$$T^{\mu\lambda} = \rho u^{\mu}u^{\lambda} + ph^{\mu\lambda} , \qquad (2.4)$$

where $h^{\mu\lambda}=g^{\mu\lambda}+u^{\mu}u^{\lambda}$ is the projection tensor, ρ is the energy density and p is the equilibrium pressure. Any other term in the above $T^{\mu\lambda}$ represents a non-equilibrium process.

In a spatially homogeneous and isotropic spacetime the only irreversible process that can appear in the energy-momentum tensor is the bulk viscosity π so out of equilibrium we have

$$T^{\mu\lambda} = \rho u^{\mu}u^{\lambda} + (p+\pi)h^{\mu\lambda}. \qquad (2.5)$$

In the background (2.1) the energy conservation law and charge conservation law are given by:

$$\rho' + (\rho + p + \pi)\theta = 0 \qquad (2.6)$$

and

$$q' + q\theta = 0 \qquad (2.7)$$

where $\theta=3R'/R$ is the scalar expansion and ' means time derivative.

In such a background the entropy flux vector S^{μ} of the first order thermodynamics is the same function of the equilibrium theory, but in the second order thermodynamics S^μ has its functional form changed and in general depends on the non-equilibrium terms that appears in the energy-momentum tensor[24,27]. In the present case S^μ is given by

$$S^\mu = (q\sigma - \frac{\alpha\pi^2}{2T})\, u^\mu \qquad (2.8)$$

where σ is the equilibrium specific entropy (entropy per unit of charge), T is the temperature and α is a coefficient to be determined.

The specific entropy obeys the equilibrium Gibbs law[28]

$$T\, d\sigma = d(\rho/q) + p\, d(1/q). \qquad (2.9)$$

Now we are ready to obtain the phenomenological relations which ensure the growth of the entropy. Taking the divergence of S^μ given by (2.8) we have, in the second order approximation,

$$S^\mu{}_{;\mu} = -\frac{\pi}{T}(\theta + \alpha\pi'). \qquad (2.10)$$

We used the equations (2.6,7,9) and the supposition that the time derivative of the coefficient α is of first order. Observe that π, θ and T' are of first order because they vanish in equilibrium.

The second law of thermodynamics ($S^{\mu}{}_{;\mu} \geq 0$) will be satisfied if π is given by (see eq. (2.10))

$$\pi = -\zeta(\theta + \alpha\pi'), \qquad (2.11)$$

where the positive coefficient of proportionality ζ is the bulk viscosity coefficient present in the first order theory, and $\alpha = \tau/\zeta$, where τ is the bulk relaxation time. Note that in the second order approximation a transient term appears which can guarantee the causality of the theory.

We suppose that soon after the GUT phase transition holds the usual equation of state

$$p = (\gamma(T) - 1)\,\rho \qquad (2.12)$$

where γ is the "adiabatic index". In fact, in the thermodynamic derivation of the above equation the γ parameter is held constant[29]. However, for a more realistic treatment, the γ index could be a slowly varying temperature function[30,31]. Moreover, the choice of the above equation of state will permit us to use the simple fluid approximation to the material content which has more than one component.

Let us now obtain the conditions for inflation. We will use the Gibbs-Duhen equation

$$dp = q\sigma \, dT + q \, d\mu, \tag{2.13}$$

where

$$\mu = (\rho + p)/q - T\,\sigma \tag{2.14}$$

is the relativistic bosonic chemical potential. In order to guarantee the non negativity of the particle and antiparticle number μ must satisfy $|\mu| \leq m$, where m is the mass of the bosons[14]. We will also use the Einstein's field equation:

$$H^2 = \frac{R'^2}{R^2} = \frac{8\pi\rho}{3m_{pl}^2}, \tag{2.15}$$

where $m_{pl}=1.22\times10^{19}$ GeV is the Planck mass.

Now, using (2.12-15) we have:

$$\frac{(1-y)}{y}\frac{H'}{H} = \frac{1}{2}\left[\left(\frac{y'}{y} - \frac{T'}{T}\right) - \frac{\mu q}{y\rho}\left(\frac{\mu'}{\mu} - \frac{T'}{T}\right)\right] \tag{2.16}$$

where $y'=T'(dy/dT)$ and $\mu'=T'(\partial\mu/\partial T)_q+q'(\partial\mu/\partial q)_T$. The above equation generalizes the expression (13) of ref.[4]. We remark that in the cited paper the chemical potential was incorrectly associated to the massless to massive boson transition. In fact, the chemical potential is a state variable associated with the existence of a conserved charge in the system. It vanishes if the net charge is zero.

From (2.16) we see that the condition for exponential inflation ($H'=0$) is given by:

$$\left(\frac{y'}{y}-\frac{T'}{T}\right)=\frac{\mu q}{y\rho}\left(\frac{\mu'}{\mu}-\frac{T'}{T}\right). \qquad (2.17)$$

We must impose

$$\frac{\mu q}{y\rho}\leq 1 \qquad (2.18)$$

to ensure the non negativity of the entropy (see eq. (2.14)). Observe that for an ideal Bose gas of particles of mass m this relation is automatically satisfied because $\rho\geq mq\geq\mu q$ and $y\geq 1$.

From (2.17) we see that there are several possibilities for having exponential inflation: When $\mu=0$ we have the condition $\ln(y/T)$=constant, that was obtained in ref.[4]. Seemingly this case is not physically accomplished, as for $\mu=0$ the system will not be dense.

When $\mu\neq 0$ the following cases are readily obtained:

(a) $y=\alpha T$ and $\mu=\beta T$ where α and β are constants.

(b) μ=const and

$$\frac{y'}{y}=\left(1-\frac{\mu q}{y\rho}\right)\frac{T'}{T}, \qquad (2.19)$$

or equivalently $dy/dT = q\sigma/\rho$.

(c) y=const and

$$\frac{\mu'}{\mu} = \left(1 - \frac{y\rho}{\mu q}\right) \frac{T'}{T} \tag{2.20}$$

or $d\mu/dT = -\sigma$.

The cases a) and b) can represent a fluid that changes its type as the temperature vary. For example, while the bosons acquire mass during the GUT phase transition, the equation of state of the fluid could change from radiation-like toward dust-like and after the bosons decay return to radiation-like again. In case c), during inflation, the chemical potential increases as the temperature decreases.

On the other hand, from (2.14) the condition for generalized inflation ($R''/R>0$ or $H > -H'/H > 0$) now takes the following form:

$$\frac{R'}{R} > \frac{y}{2(y-1)} \left[\left(\frac{y'}{y} - \frac{T'}{T}\right) - \frac{\mu q}{y\rho} \left(\frac{\mu'}{\mu} - \frac{T'}{T}\right) \right] > 0 \ . \tag{2.21}$$

It can also be shown that the condition (2.17) for exponential inflation does not change if one includes a cosmological constant Λ in the Einstein's equation. However, the condition for generalized inflation (2.21) changes to

$$\frac{R'}{R} > \frac{y}{2(y-1)} \left(1 - \frac{\Lambda}{3HR}\right) \left[\left(\frac{y'}{y} - \frac{T'}{T}\right) - \frac{\mu q}{y\rho} \left(\frac{\mu'}{\mu} - \frac{T'}{T}\right) \right] > 0 \ . \tag{2.22}$$

III. THE BULK VISCOSITY COEFFICIENT, ENTROPY PRODUCTION AND THE DURATION OF INFLATION

Let us obtain the bulk viscosity coefficient in the exponential inflation case. Substituting $\rho'=0$ in (2.6) and using (2.12) we have

$$\pi = -y\rho. \qquad (3.1)$$

Using the above equation in the constitutive equation (2.11) we obtain

$$\zeta = \zeta_{qs}(1 + \frac{\gamma y'}{y}) , \qquad (3.2)$$

where $\zeta_{qs}=y\rho/3H$ is the bulk viscosity coefficient of the quasi-stationary theory[32]. Note that except for the case c), there is a transient term in addition to the expression of the quasi-stationary theory. We can also express ζ as a function of temperature. For example, for the case a) we have

$$\zeta \propto T\ (1 + \frac{\gamma T'}{T}) . \qquad (3.3)$$

In the present model the temperature variation is not determined by the dynamical equations. In fact, the equations of motion are fully satisfied by the condition $\rho'=0$ and by the relation (3.1) in the exponential inflation case, while the variation of the temperature depends on the details of the transition between the phase where the bosons are massless and the phase after when they acquire mass.

Let us analyze the entropy production during the inflationary period. In the nonstationary thermodynamic theory the physical entropy density is given by (2.8) ($s=-S^{\mu}u_{\mu}$), so

$$s = q\sigma - \frac{\tau\pi^2}{2\xi T} . \qquad (3.4)$$

Using the equations (2.14), (3.1-2) and the charge conservation law, we obtain from (3.4) that:

$$s = \frac{\gamma\rho}{T}(1 - 3\tau H/2) - \frac{\mu q_0}{T}\exp(-3H(t-t_0)) \qquad (3.5)$$

where q_0 is the charge density in the beginning of the inflationary period ($t=t_0$). If $\tau=\mu=0$ the result of the ref.[4] is reobtained, namely $s=\gamma\rho/T$=constant. If $\tau' \leq 0$ then the entropy density is an increasing function of time for all cases a) to c). For the cases a) and b) the last term of (3.5) becomes negligible for $\Delta t > H^{-1}$ (see the relation (2.18)) while the duration of inflation is $\gg H^{-1}$. The second order thermodynamic works well when the second term in (3.4) is much smaller than the first one, that happens when $\tau H \ll 1$. The limit of validity of the second order theory occurs when the terms are of the same order, that is $\tau H \approx 1$.

In the last case the entropy density would be small and the duration of inflation should be too long in order to solve the flatness problem.

It is easy to estimate the duration of inflation in order to solve the entropy problem in case a) when $\tau' \approx 0$. As $R = R_0 \exp[H(t-t_0)]$ then

$$\Delta t \approx \frac{1}{3H} \ln\left(\frac{S_f}{sR_0{}^3}\right) \qquad (3.6)$$

where S_f is the entropy at the end of inflation that is $\approx 10^{87}$. Using $H \approx 10^{10}$ GeV, $R_0 \approx 10^{-12}$ GeV^{-1} and $\gamma\rho/T \approx 10^{42}$ GeV^3 follows that

$$\Delta t \approx 10^{-10} \ln \frac{10^{81}}{(1-3\tau H/2)} \mathrm{GeV}^{-1} . \qquad (3.7)$$

For $\tau = 0$ we have $\Delta t \approx 10^{-32}$ seg. Observe that as s is nearly constant, the universe radius would increase by a factor of 10^{28}, that is enough to solve the horizon, homogeneity and isotropy puzzles.

In the present model, the entropy generation is concomitant with the inflation, differently from the vacuum-pressure-driven inflationary models[33-35], where the entropy is generated after inflation by a highly non adiabatic process. During the "slow rollover phase" of the new inflationary universe, the temperature should decrease 10^{28} times in order to maintain the radiation entropy constant. Here, as the bulk viscosity reheats continuously the medium, the temperature need not to decrease so drastically.

As remarked in the introduction the material content remains dense just while the charge density is $\gg T^3$. Inflation dilutes any charge excess by a factor $\gtrsim 10^{84}$. So, admitting that at the end of the inflationary phase the system has diluted, the charge density will be $\gtrsim T_f{}^3$ (we are not taking into account the possible contributions of monopoles, cosmic strings, etc., to enhance the number density). Then, if the final temperature is not too low the initial value of q will be very high and could turn the model meaningless, since ρ could be greater than ρ_{Planck}. This problem arise due the assumption that the gauge bosons acquire mass $\approx 10^{14}$ GeV instantaneously and simultaneously. It can be avoided if at least one of these conditions is relaxed during the phase transition.

We now use the limit of validity ($m \approx T$) of equation (1.2) to rough estimate the final decondensation critical temperature in the isothermal case. It is easy to see that if $q_f \approx T^3$ then $T_c \approx T$. Since the beginning of inflation T_c is a decreasing function of time then this estimate shows us that the decondensation process will occur together with the dilution process.

IV. CONCLUSION

In this paper we have considered the consequences of a net bosonic charge present in the universe during the inflationary period. We have showed that even in the realm of asymptotic freedom, the existence of a charge can avoid the validity of the dilute gas approximation at the GUT epoch permitting the bulk viscosity to generate inflation. We have discussed the presence of a Bose-Einstein condensation during the GUT phase transition and pointed out that, due to the universe expansion, it could occur another phase transition representing the decondensation process.

The conditions for inflation have been obtained in both exponential and generalized case. The expression for the bulk viscosity coefficient that was previously derived has been extended by using causal thermodynamic theory. We have also showed that in the context of this theory the duration of inflation is longer. In the present model inflation ends by dilution. As the universe expands the charge density decreases and the ratio λ/L becomes $\gg 1$. In this case bulk viscosity can not drive inflation anymore.

In the course of our investigation some simplifications were performed. Firstly, we have used a linear approximation in the thermodynamical phenomenological laws. Secondly, we have used only one chemical potential, however more than one charge could exist. Thirdly, we have described the primordial cosmic plasma as a simple fluid with the equation of state given by (2.12). In fact, we think that a more realistic treatment is necessary for a full description of the universe at those eras. Further investigations in this direction are being accomplished.

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