THE QUASI-STEADY STATE COSMOLOGY

F. HOYLE
102 Admirals Walk, West Cliff Road, Dorset, BH2 5HF Bournemouth, United Kingdom

G. BURBIDGE
University of California, San Diego, Department of Physics and Center for Astrophysics & Space Sciences, 9500 Gilman Drive, La Jolla, California 92093-0111

AND

J.V. NARLIKAR
Inter-University Centre for Astronomy & Astrophysics, Poona University Campus, Pune 411 007, India

1. Introduction

At IAU Symposium No. 168 the presentation was divided into two parts:
I. Theoretical Foundations and
II. Observational Facts and Consequences

Nearly all of the work presented here is contained in four papers published by Hoyle, Burbidge and Narlikar (1993; 1994a,b,c) which will be abbreviated in the text to HBN 1993, HBN 1994a, HBN 1994b and HBN 1994c.

In this presentation Section 2 is devoted to the basic theory and in Section 3 we describe the observations and the way that we interpret them using the theory. Also in Section 3 we discuss various predictions relating to the theory.

2. Theoretical Foundations

To begin with we show with the help of a model how the problems of space-time singularity and violation of the energy momentum conservation law that are present in the standard cosmology can be avoided by introducing

M. Kafatos and Y. Kondo (eds.), Examining the Big Bang and Diffuse Background Radiations, 329–368.

a scalar field minimally coupled to gravity and having its sources in events where matter is created.

We than show that matter creation preferentially occurs near collapsed massive objects and the scalar field created at such mini-creation events has a feedback on spacetime geometry causing the universe to have a steady expansion as in the de Sitter model but with periodic phases of expansion and contraction superposed on it.

The parameters of our model can be empirically fixed in relation to the cosmological observations thus providing tests of the theory. In the second part the observational aspects are dealt with.

Next we argue that the model arises from a deeper theory which is Machian in origin with the inertia of a particle determined by the rest of the particles in the universe in a long range conformably invariant scalar interaction. The characteristic mass of a particle created is then the Planck mass. The Planck particle decays quickly to baryons. The inertial effects produced by the Planck particles during their brief existence generate the scalar field of the toy model while the inertial effects of the stable baryonic particles give the more familiar Einstein equations of relativity.

Finally we show that extending the theory to the most general conformably invariant form automatically leads to the cosmological constant whose sign and magnitude are of the right cosmological order.

We begin with the tentative definition that cosmology refers to a study of those aspects of the universe for which spatial isotropy and homogeneity can be used, with the spacetime metric taking the form

$$ds^2 \quad = \quad dt^2 - S^2(t)\left[\frac{dr^2}{1-kr^2} + r^2(d\theta^2 + sin^2\theta d\phi^2)\right] \tag{1}$$

in terms of coordinates t, r, θ, ϕ with $r = 0$ at the observer. The topological constant k in this so-called Robertson–Walker form can be shown to be 0 or ± 1. The "particles" to which (1) applies are thought of as galaxies or clusters of galaxies, each "particle" having spatial coordinates r, θ, ϕ independent of the universal time t. They form what is often referred to as the Hubble flow.

Big-Bang cosmology in all its forms is obtained from the equations of general relativity,

$$R_{ik} \quad - \quad \frac{1}{2}g_{ik}R + \lambda g_{ik} = -8\pi G T_{ik}, \tag{2}$$

which follow from the variation of an action formula

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$$\mathcal{A} = \frac{1}{16\pi G}\int_V (R+2\lambda)\sqrt{-g}d^4x + \int_V \mathcal{L}_{phys}(X)\sqrt{-g}d^4x \qquad (3)$$

with respect to a general Riemannian metric

$$ds^2 = g_{ik}dx^i dx^k \qquad (4)$$

within a general spacetime volume V. The physical Lagrangian $\mathcal{L}_{phys}(X)$ generates the energy-momentum tensor T_{ik} in this variation of g_{ik}. In the standard big-bang cosmology the physical Lagrangian includes only particles and the electromagnetic field, whereas in inflationary forms of big-bang cosmology a scalar field is also considered to be added to $\mathcal{L}_{phys}$. This is done in various ways, being severally advocated by different authors (see Narlikar and Padmanabhan 1991 for a review).

The initial conditions assumed in the standard model are :

(i) The universe was sufficiently homogeneous and isotropic at the outset for the metric (1) to be used immediately over a range of the r-coordinate of relevance to presentday observation,
(ii) $k = 0$,
(iii) $\lambda = 0$,
(iv) The initial balance of radiation and baryonic matter was such that the light elements D, ^{3}He, ^{4}He, ^{7}Li were synthesised in the early universe in the following relative abundances to hydrogen

$$\frac{D}{H} \simeq \frac{^3He}{H} \simeq 2\times 10^{-5}, \frac{^4He}{H} \simeq 0.235, \frac{^7Li}{H} \simeq 10^{-10}.$$

From detailed calculations these abundances can be shown to require

$$\rho_{\text{baryon}} \simeq 10^{-32}T^3 \text{ gcm}^{-3}, \qquad (5)$$

the radiation temperature being in degrees kelvin. (Gamow 1946, Alpher, et al 1950, 1953, Hoyle and Tayler 1964).

There is a fundamental problem here. The instant t = 0 is the so-called spacetime singularity at which the field equations (2) break down. This is identified with the big bang epoch. All matter that we see in the universe (as well as radiation) is supposed to be given as an initial condition at $t = \epsilon > 0$. The initial instant ϵ can be taken arbitrarily close to $t = 0$ but not identified with it. Thus the action principle (3) itself gets restricted in validity since the singular epoch must be excluded from it too.

Conceptually this is an exceptional step to take. In theoretical physics the basic laws or principles like the action principle are considered superior to the specific solutions based on them. Yet here we seek to restrict the validity of (2) and (3) because the solution so warrants it! There is thus a clear indication here of an inconsistency of the overall framework.

The other problems of the standard big bang model often referred to as the horizon and flatness problems also relate to the above initial conditions assumed at $t = \epsilon > 0$. While the need for such far reaching assumptions as (i) to (iv) has always prompted a measure of unease they were widely accepted for a decade and a half, and are indeed still fully accepted by the more orthodox supporters of the standard model. Others, however, welcomed the inflationary idea of including a scalar field in the physical lagrangian that initially dominated both matter and other fields and which varied adiabatically in such a way as to give

$$\frac{\dot{S}^2}{S^2} = C, \quad S(t) = S(0) \exp \sqrt{C}t, \tag{6}$$

with C a constant. The solution (6) is considered to apply from $t = \epsilon > 0$, where $\epsilon \sim 10^{-36}$s to a value of t large compared to $1/\sqrt{C}$. It greatly reduces the range of the r-coordinate over which (i) is needed and it effectively removes the k-term from (1). It also removes any initial contributions from matter and radiation, but these are considered to be reasserted through a physical transition of the scalar field, which jumps the solution (6) to

$$\dot{S}^2 \ = \ \frac{A}{S}, S \simeq \left(\frac{9}{4}A\right)^{\frac{1}{3}} t^{\frac{2}{3}}, \tag{7}$$

which is the so-called closure model with matter just having sufficient expansion to reach a state of infinite dispersal, a condition that is considered most favorable for the eventual formation of stars and galaxies.

A major problem associated with inflation is how to effectively eliminate the cosmological constant. The value of this constant which gave the exponential solution (6) above must reduce to zero or, if the cosmological observations so demand, become as small as 10^{-108} of its initial value.

Any theoretical trick invoked to achieve this has a contrived appearance (Weinberg 1989).

The papers referred to earlier (HBN 1993, 1994abc) show how we have developed an alternative scenario.

The action principle (3) has a second term which is supposed to include physical contributions other than gravity. A close parallel exists between the scalar field used for inflation and the scalar field used earlier by Hoyle and Narlikar (1963) for obtaining the steady state model from Einstein's field equations. To begin with we will use the 1963 formalism as a "toy model" for describing creation of matter without violating the law of conservation of energy-momentum and without encountering spacetime singularity.

Thus the classical Hilbert action leading to the Einstein equations is modified by the inclusion of a scalar field C whose derivatives with respect to the spacetime coordinates x^i are denoted by C_i. The action is given by

$$\begin{aligned} \mathcal{A} \;=\; & -\sum_a \int_{\Gamma_a} m_a ds_a + \int_V \frac{1}{16\pi G} R\sqrt{-g}d^4x - \frac{1}{2}f\int_V C_i C^i \sqrt{-g}d^4x \\ & + \sum_a \int_{\Gamma_a} C_i da^i \end{aligned} \tag{8}$$

where C is a scalar field and $C_i = \partial C/\partial x^i$. f is a coupling constant. The last term of (8) is manifestly path-independent and so, at first sight it appears to contribute no new physics. The first impression, however, turns out to be false if we admit the existence of broken worldlines. For, if particles $a, b, \ldots$ are created at world points $A_0, B_0, \ldots$ respectively, then the last term of (8) contributes a non-trivial sum

$$-\{C(A_0) + C(B_0) + \ldots\}$$

to $\mathcal{A}$.

Thus, if the worldline of particle a begins at point A_0, then the variation of $\mathcal{A}$ with respect to that worldline gives

$$m_a \frac{da^i}{ds_a} \;=\; g^{ik} C_k \tag{9}$$

at A_0. In other words, the C-field balances the energy-momentum of the created particle.

The field equations likewise get modified to

$$R_{ik} - \frac{1}{2}g_{ik}R = -8\pi G\left[\underset{m}{T_{ik}} + \underset{c}{T_{ik}}\right] \tag{10}$$

where

$$\underset{c}{T_{ik}} = -f\left\{C_iC_k - \frac{1}{2}g_{ik}C^lC_l\right\}. \tag{11}$$

Thus the energy conservation law is

$$\underset{m}{T^{ik}}{}_{;k} = -\underset{c}{T^{ik}}{}_{;k} = fC^iC^k{}_{;k}. \tag{12}$$

That is, matter creation via a nonzero left hand side of (12) is possible while conserving the overall energy and momentum. The C-field tensor has negative stresses which lead to the expansion of spacetime, as in the case of inflation.

From (9) we therefore get a necessary condition for creation as

$$C_iC^i = m_a^2; \tag{13}$$

this is the 'creation threshold' which must be crossed for particle creation. How this can happen near a massive object, can be seen from the following simple example.

The Schwarzschild solution for a massive object M of radius $R > 2GM/c^2$ is

$$ds^2 = dt^2\left(1 - \frac{2GM}{r}\right) - \frac{dr^2}{1 - \frac{2GM}{r}} - r^2(d\theta^2 + \sin^2\theta d\phi^2), \tag{14}$$

for $r \geq R$. Now if the C-field does not seriously change the geometry, we would have at $r \gg R$,

$$\dot{C} \approx m, \qquad C' \equiv \frac{\partial C}{\partial r} \cong 0. \tag{15}$$

If we continue this solution closer to $r \approx R$, we find that

$$C^i C_i \equiv \left(1 - \frac{2GM}{r}\right)^{-1} m^2. \qquad (16)$$

In other words $C_i C^i$ increases towards the object and can become arbitrarily large if $r \approx 2GM$. So it is possible for the creation threshold to be reached *near* a massive collapsed object even if $C_i C^i$ is *below* the threshold far away from the object. In this way massive collapsed objects can provide new sites for matter creation. Further, because of the negative stresses the created matter is expelled outwards from the site while the C-field quanta escape with the speed of light. Thus, instead of a single big bang event of creation, we have mini-creation events near collapsed massive objects.

Since the C-field is a global cosmological field, we expect the creation phenomenon to be globally cophased. Thus, there will be phases when the creation activity is large, leading to the generation of the C-field strength in large quantities. However, the C-field growth because of its large negative stresses leads to a rapid expansion of the universe and a consequent drop in its background strength. When that happens creation is reduced and takes place only near the most collapsed massive objects thus leading to a drop in the intensity of the C-field. The reduction in C-field slows down the expansion, even leading to local contraction and so to a build-up of the C-field strength. And so on!

We can describe this up and down type of activity as an oscillatory solution superposed on a steadily expanding de Sitter type solution of the field equations as follows. For the Robertson–Walker line element the equations (10)–(12) give

$$3\frac{\dot{S}^2 + kc^2}{S^2} = 8\pi G(\rho - \frac{1}{2} f \dot{C}^2), \qquad (17)$$

$$2\frac{\ddot{S}^2}{S} + \frac{\dot{S}^2 + kc^2}{S^2} = 4\pi G f \dot{C}^2, \qquad (18)$$

where $S(t)$ is the scale factor and k the curvature parameter $(= 0, \pm 1)$. The cosmic time is given by t. These equations have a deSitter type solution given by

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$$S \propto \exp(t/P), k = 0, \qquad \dot{C} = \text{constant}, \qquad \rho \;=\; \text{constant} \tag{19}$$

The oscillatory solution is given by

$$k = +1, \qquad \dot{C} \propto 1/S^3, \qquad \rho \propto 1/S^3. \tag{20}$$

Thus (17) becomes, in the latter case

$$\dot{S}^2 \;=\; -c^2 + \frac{A}{S} - \frac{B}{S^4}, \quad A, B = \text{constant}. \tag{21}$$

Here the oscillatory cycle will typically have a period $Q \ll P$.

Although the exact solution of (21) will be difficult to obtain, we can use the following approximate solution of (19) and (20) to describe the short-term and long-term cosmological behaviour :

$$S(t) \;=\; \exp\left(\frac{t}{P}\right)\left\{1 + \alpha\cos\frac{2\pi t}{Q}\right\}. \tag{22}$$

Note that the universe has a long term secular expanding trend, but because $|\alpha| < 1$, it also executes non-singular oscillations around it. For this reason this model has been called "quasi-steady state cosmology". We can determine α and our present epoch $t = t_0$ by the observations of the present state of the universe. Thus an acceptable set of parameters is

$$\alpha \;=\; 0.75,\; t_0 = 0.85Q,\; Q = 4 \times 10^{10} yr.,\; P = 20Q. \tag{23}$$

We shall show how their values relate to the observed parameters in Part II.

But now we discuss the theory underlying the physics of matter creation.

An important property of physical theories is scale invariance or conformal invariance. Maxwell's equations and the Dirac equation for a massless particle are conformably invariant but general relativity is not. If, however, the inertial mass transforms inversely as the length scale in conformal transformation then the Dirac equation for a massive fermion as well as classical and quantum electrodynamics will become conformably invariant. Can general relativity be suitably reformulated to be conformably invariant? We indicate the steps towards this goal since they naturally lead to a comprehensive theory of matter creation that encompasses our model.

It is necessary to begin by finding an action $\mathcal{A}$ that is unaffected in its value by a scale transformation. The second term on the right-hand side of (3) can be made to satisfy this requirement. For a set of particles $a, b, \ldots$ of masses $m_a, m_b, \ldots$ the form of $\mathcal{L}_{\text{phys}}$ usually considered in gravitational theory is

$$\sum_{a,b,\ldots} \int \frac{\delta_4(X,A)}{\sqrt{-g(A)}} m_a(A)da, \tag{24}$$

where the possibility of the particle masses varying with the spacetime position requires the mass $m_a(A)$ of particle a to vary with the point A on its path, and similarly for the other particles. Hence the second term on the right-hand side of (3) is $-\sum_a \int m_a(A)da$.
With $da^* = \Omega da$ and $m_a^* = \Omega^{-1} m_a$ it is clear that (25) is invariant with respect to a conformal (scale) transformation.

$$\Box_X M(X) + \frac{1}{6} R M(X) = \sum_a \int \frac{\delta_4(X,A)}{\sqrt{-g(A)}} da. \tag{25}$$

The possibility of particle masses varying with spacetime coordinates arises most naturally in a Machian approach. Here the property of inertia is not entirely intrinsic to a particle but is also related to its presence in a non-empty universe. A quantitative description of this idea that we will follow here is based on an early work by two of us (Hoyle and Narlikar 1964). In this inertia is expressed as a scalar conformably invariant long range interaction between particles.

To begin with choose a scalar mass field $M(X)$ to satisfy

$$-\sum_a \int m_a(A)da. \tag{26}$$

Equation (26) has both advanced and retarded solutions. We particularize an advanced solution $M^{\text{adv}}(X)$ and a retarded solution $M^{\text{ret}}(X)$ in the following way. $M^{\text{ret}}(X)$ is to be the so-called fundamental solution in the flat spacetime limit (Courant and Hilbert, 1962). This removes for $M^{\text{ret}}(X)$

the ambiguity that would obviously arise from the homogeneous wave equation. The corresponding ambiguity for $M^{\rm adv}(X)$ is removed by the physical requirement that fields without sources are to be zero. Since

$$\left[M^{\rm adv} - M^{\rm ret}\right] + \frac{1}{6}R[M^{\rm adv} - M^{\rm ret}] \;=\; 0, \tag{27}$$

the immediate consequences of this boundary condition is that $M^{\rm adv} - M^{\rm ret}$, being without sources, must be zero, so that

$$M^{\rm adv}(X) \;=\; M^{\rm ret}(X) = M(X) \text{ say}. \tag{28}$$

The gravitational equations are now obtained by putting

$$m_a(A) \;=\; M(A), \quad m_b(B) = M(B), \ldots \tag{29}$$

It can also be shown that in a conformal transformation the mass field $M(X)$ transforms as

$$M^*(X) \;=\; \Omega^{-1}(X)M(X), \tag{30}$$

a result that follows from the form of the wave equation (10) (c.f. Hoyle and Narlikar, 1974, 111). The outcome (*loc. cit.*, 112 *et seq*) is

$$\begin{aligned} K(R_{ik} - \frac{1}{2}g_{ik}R) \;&=\; -T_{ik} + M_i M_k - \frac{1}{2}g_{ik}g^{pq}M_p M_q \\ &\quad + G_{ik}K - K_{;ik}, \end{aligned} \tag{31}$$

where

$$K \;=\; \frac{1}{6}M^2. \tag{32}$$

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These gravitational equations are scale invariant. It may seem curious that from a similar beginning, (24) for the action rather than (3), the outcome is more complicated, but this seems to be a characteristic of the physical laws. As the laws are improved they become simpler and more elegant in their initial statement but more complicated in their consequences.

Now make the scale change

$$\Omega(X) \quad = \quad M(X)/\tilde{m}_0, \tag{33}$$

where $\tilde{m}_0$ is a constant with the dimensionality of $M(X)$. After the scale change the particle masses simply become $\tilde{m}_0$ everywhere and in terms of transformed masses the derivative terms drop out of the gravitational equations. And defining the gravitational constant G by

$$G \quad = \quad \frac{3}{4\pi\tilde{m}_0^2}, \tag{34}$$

the equations (31) take the form of general relativity

$$R_{ik} - \frac{1}{2}g_{ik}R \quad = \quad -8\pi G T_{ik}. \tag{35}$$

It now becomes clear why the equations of general relativity are not scale invariant. These are the special form to which the scale invariant equations (31) reduce with respect to a particular scale, namely that in which particle masses are everywhere the same.

It is also clear that the transition from (31) to (35) is justified provided $\Omega(X) \neq 0$ or $\Omega(X) \neq \infty$. For example, if $M(X) = 0$ on a spacelike hypersurface the above conformal transformation breaks down. It is because of the existence of such time sections that the use of (35) leads to the (unphysical) conclusion of a spacetime singularity. It was shown (Hoyle and Narlikar 1974, Kembhavi 1979) that the various spacetime singularities like that in the big bang or in a black hole collapse arise because of the illegitimate use of (35) in place of (31).

It is easily seen from the wave equation (26) that $M(X)$ has dimensionality (length)$^{-1}$, and so has $\tilde{m}_0$. Units are frequently used in particle physics for which both the speed of light c and Planck's constant $\hbar$ are unity and in these units mass has dimensionality (length)$^{-1}$. If we suppose these units apply to the above discussion then from (34)

$$\tilde{m}_0 \;=\; (3/4\pi G)^{1/2}, \tag{36}$$

which with $c = \hbar = 1$ is the mass of the Planck particle. This suggests that in a gravitational theory without other physical interactions the particles must be of mass (36), which in ordinary practical units is about 10^{-5} gram, the empirically determined value of G being used. This conclusion can be supported at greater length [See HBN 1994c]. We next consider what happens when the Planck mass decays into the much more stable baryons.

A typical Planck particle a exists from A_0 to $A_0 + \delta A_0$, in the neighborhood of which it decays into n stable baryonic secondaries, $n \simeq 6.10^{18}$, denoted by $a_1, a_2, \ldots a_n$. Each such secondary contributes a mass field $m^{(a_r)}(X)$, say, which is the fundamental solution of the wave equation

$$m^{(a_r)} + \frac{1}{6} R m^{(a_r)} \;=\; \frac{1}{n} \int_{\sim A_0 + \delta A_0} \frac{\delta_4(X, A)}{\sqrt{-g(A)}} da, \tag{37}$$

while the brief existence of a contributes $c^{(a)}(X)$, say, which satisfies

$$c^{(a)} + \frac{1}{6} R c^{(a)} \;=\; \int_{A_0}^{A_0 + \delta A_0} \frac{\delta_4(X, A)}{\sqrt{-g(A)}} da. \tag{38}$$

Summing $c^{(a)}$ with respect to $a, b, \ldots$ gives

$$c(X) \;=\; \sum_a c^{(a)}(X), \tag{39}$$

the contribution to the total mass $M(X)$ from the Planck particles during their brief existence, while

$$\sum_a \sum_{r=1}^{n} m^{(a_r)}(X) \;=\; m(X) \tag{40}$$

gives the contribution of the stable particles.

Although $c(X)$ makes a contribution to the total mass function

$$M(X) = c(X) + m(X) \tag{41}$$

that is generally small compared to $M(X)$, there is the difference that, whereas $m(X)$ is an essentially smooth field, $c(X)$ contains small exceedingly rapid fluctuations and so can contribute significantly to the derivatives of $c(X)$. The contribution to $c(X)$ from Planck particles a, for example, is largely contained between two light cones, one from A_0, the other from $A_0 + \delta A_0$. Along a timelike line cutting these two cones the contribution to $c(X)$ rises from zero as the line crosses the light cone from A_0, attains some maximum value and then falls back effectively to zero as the line crosses the second light cone from $A_0 + \delta A_0$. The time derivative of $c^{(a)}(X)$ therefore involves the reciprocal of the time difference between the two light cones. This reciprocal cancels the short duration of the source term on the right-hand side of (40). The factor in question is of the order of the decay time τ of the Planck particles, $\sim 10^{-43}$ seconds. No matter how small τ may be the reduction in the source strength of $c^{(a)}(X)$ is recovered in the derivatives of $c^{(a)}(X)$, which therefore cannot be omitted from the gravitational equations.

The derivatives of $c^{(a)}(X), c^{(b)}(X), \ldots$ can as well be negative as positive, so that in averaging many Planck particles, linear terms in the derivatives do disappear. It is therefore not hard to show that after such an averaging the gravitational equations become

$$R_{ik} - \frac{1}{2} g_{ik} R = \frac{6}{m^2} \left[-T_{ik} + \frac{1}{6}(g_{ik} m^2 - m^2_{;ik}) + (m_i m_k - \frac{1}{2} g_{ik} m_l m^l) + \frac{2}{3}\left(c_i c_k - \frac{1}{4} g_{ik} c_l c^l\right)\right]. \tag{42}$$

Since the same wave equation is being used for $c(X)$ as for $m(X)$, the theory remains scale invariant. A scale change can therefore be introduced that reduces $M(X) = m(X) + c(X)$ to a constant, or one that reduces $m(X)$ to a constant. Only that which reduces $m(X)$ to a constant, viz

$$\Omega = \frac{m(X)}{m_0} \tag{43}$$

has the virtue of not introducing small very rapidly varying ripples into the metric tensor. Although small in amplitude such ripples produce non-

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negligible contributions to the derivatives of the metric tensor, causing difficulties in the evaluation of the Riemann tensor, and so are better avoided. Simplifying with (43) does not bring in this difficulty, which is why separating of the main smooth part of $M(X)$ in (41) now proves an advantage, with the gravitational equations simplifying to

$$8\pi G = \frac{6}{m_0^2}, \quad m_0 \text{ a constant,} \tag{44}$$

$$R_{ik} - \frac{1}{2}g_{ik}R = -8\pi G[T_{ik} - \frac{2}{3}(c_i c_k - \frac{1}{4}g_{ik}c_l c^l)]. \tag{45}$$

Using the metric (1) with $k = 0$ the dynamical equations for the scale factor $S(t)$ are

$$\frac{2\ddot{S}}{S} + \frac{\dot{S}^2}{S^2} = \frac{4\pi}{3}G\bar{\dot{c}^2}, \tag{46}$$

$$\frac{3\dot{S}^2}{S^2} = 8\pi G\left(\bar{\rho} - \frac{1}{2}\bar{\dot{c}^2}\right), \tag{47}$$

with $\bar{\rho}$ the average particle mass density and $\bar{\dot{c}^2}$ being the average value of $\dot{c}^2$, the average value of terms linear in c and of $\ddot{c}$ being zero. It is easily shown from (46) and (47) that

$$\frac{\partial\bar{\rho}}{\partial t} + \frac{3\dot{S}}{S}\bar{\rho} = \frac{1}{2}\left(\frac{\partial\bar{\dot{c}^2}}{\partial t} + \frac{4\dot{S}}{S}\bar{\dot{c}^2}\right). \tag{48}$$

If at a particular time there is no creation of matter then at that time the left-hand side of (48) is zero with $\bar{\rho} \propto S^{-3}$. And with the right-hand side also zero at that time $\bar{\dot{c}^2} \propto S^{-4}$. The sign of the $\bar{\dot{c}^2}$ term in (46) is that of a negative pressure, a characteristic of the fields introduced into inflationary cosmological models. The concept of Planck particles forces the appearance of a negative pressure. In effect the positive energy of created particles is compensated by the sign of the $\bar{\dot{c}^2}$ terms, which in (46) increases

$\ddot{S}/S$ and so causes the universe to expand. One can say that the universe expands because of the creation of matter. The two are connected because the divergence of the right-hand side of the gravitational equations (45) is zero.

As would be expected from this conservation property the sign of the $\dot{c}^2$ term in (47) is that of a negative energy field. Such fields have generally been avoided in physics because in flat spacetime they would produce catastrophic instabilities – creation of matter with positive energy producing a negative energy $\dot{c}^2$ term, producing more matter, producing a still larger $\dot{c}^2$ term, and so on. Here the effect is to produce explosive outbursts from regions where any such instability takes hold, through the $\dot{c}^2$ term in (46) generating a sharp increase of $\ddot{S}$. The sites of the creation of matter are thus potentially explosive. The explosive expansion of space serves to control the creation process and avoids the catastrophic cascading down the negative energy levels.

As will be discussed in II, this is in agreement with observational astrophysics which in respect to high energy activity is all of explosive outbursts, without evidence for the ingoing motions required by the popular accretion-disk theory for which there is no direct observational support. The profusion of sites where X-ray and γ-ray activity is occurring are on the present theory sites where the creation of matter is currently taking place.

A connection with our model can now be given. Writing

$$C(X) \quad = \quad \tau c(X), \tag{49}$$

where τ is the decay lifetime of the Planck particle, the action contributed by Planck particles $a, b, \ldots$,

$$-\sum_a \int_{A_0}^{A_0+\delta A_0} c(A)da \tag{50}$$

can be approximated as

$$-C(A_0) - C(B_0) - \ldots, \tag{51}$$

which form was used in the model. And the wave-equation for $C(X)$, using the same approximation, is

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$$C + \frac{1}{6}RC \;=\; \tau^{-2}\sum_a \frac{\delta_4(X, A_0)}{\sqrt{-g(A_0)}}, \tag{52}$$

which was also used in the model, except that the 1/6 RC term is included in the wave equation and previously an unknown constant f appeared in place of τ^2.

Writing $M^{(a)}(X), M^{(b)}(X), \ldots$ as the mass fields produced by the individual Planck particles $a, b, \ldots$, the total mass field

$$M(X) \;=\; \sum_a M^{(a)}(X) \tag{53}$$

satisfies the wave equation (26) when $M^{(a)}, M^{(b)}, \ldots$ satisfy

$$M^{(a)} + \frac{1}{6}RM^{(a)} \;=\; \int \frac{\delta_4(X, A)}{\sqrt{-g(A)}} da, \ldots \tag{54}$$

Scale invariance throughout requires all the mass fields to transform as

$$M^{*(a)} \;=\; M^{(a)}\Omega^{-1} \tag{55}$$

with respect to the scale change Ω, when both the left and right hand sides of every wave equation transform to its starred form multiplied by Ω^{-3}, i.e. the left hand side of (54) goes to $(M^{*(a)} + \frac{1}{6}R^*M^{*(a)})\Omega^{-3}$ and the right hand side to

$$\Omega^{-3}\int \frac{\delta_4(X, A)}{\sqrt{-g^*(A)}} da^*. \tag{56}$$

Then the factor Ω^{-3} cancels to give the appropriate invariant equation. This cancellation is evidently unaffected if, instead of (54) for the wave equation satisfied by M^a, we have

$$M^{(a)} + \frac{1}{6}RM^{(a)} + M^{(a)^3} \;=\; \int \frac{\delta_4(X, A)}{\sqrt{-g(A)}} da. \tag{57}$$

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Since the cube term transforms to $M^{*(a)^3}\Omega^{-3}$ with respect to Ω changing (54) to (57) preserves scale invariance in what appears to be its widest form. Since in other respects the laws of physics always seem to take on the widest ranging properties that are consistent with the relevant forms of invariance we might think it should also be here, in which case (57) rather than (54) is the correct wave equation for $M^{(a)}$, and similarly for $M^{(b)},\ldots$, the mass fields of the other Planck particles.

But this departure from linearity in the wave equations for the individual particles prevents a similar equation being obtained for $M = \sum_a M^{(a)}$. Nevertheless, the addition of the individual equations can be considered in a homogeneous universe to lead to an approximate wave equation for M of the form

$$M + \frac{1}{6}RM + \Lambda M^3 \;=\; \sum_a \int \frac{\delta_4(X,A)}{\sqrt{-g(A)}}da, \tag{58}$$

$$\Lambda \;=\; N^{-2}, \tag{59}$$

where N is the effective number of particles contributing to the sum $\sum_a M^{(a)}$. The latter can be considered to be determined by an Olbers-like cut-off, contributed by the portion of the universe surrounding the point X in $M(X)$ to a redshift of order unity. In the observed universe this total mass $\sim 10^{22}M_\odot$, sufficient for $\sim 2.10^{60}$ Planck particles. The actual particles are of course nucleons of which there are $\sim 10^{79}$ But if suitably aggregated they would give $\sim 2.10^{60}$ Planck particles and with this value for N

$$\Lambda \;\simeq\; 2.5\times 10^{-121}. \tag{60}$$

The next step is to notice that the wave-equation (58) would be obtained in usual field theory from $\delta\mathcal{A} = 0$ for $M \to M + \delta M$ applied to

$$\begin{aligned}\mathcal{A} \;=\; & -\frac{1}{2}\int (M_i M^i - \frac{1}{6}RM^2)\sqrt{-g d^4x} + \frac{1}{4}\Lambda \int M^4\sqrt{-g d^4 x} \\ & -\sum_a \int \frac{\delta_4(X,A)}{\sqrt{-g(A)}}M(X)da.\end{aligned} \tag{61}$$

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In the scale in which M is m_0 everywhere the derivative term in (61) vanishes and since $G = 3/4\pi m_0^2$ the term in R is the same as in (3), as are also the line integrals, requiring the remaining term to be the same gives

$$\lambda = -3\Lambda m_0^2. \tag{62}$$

Thus we have obtained not only a cosmological constant but also its magnitude, something that lies beyond the scope of the usual theory. With 2.5×10^{-121} for Λ as in (60) and with m_0 the inverse of the Compton wavelength of the Planck particle, $\sim 3.10^{32}$ cm^{-1}, (62) gives

$$\lambda \simeq -2.10^{-56}\ \mathrm{cm}^{-2}, \tag{63}$$

agreeing closely with the magnitude that has previously been assumed for λ. In the classical big bang cosmology there is no dynamical theory to relate the magnitude of λ to the density or other physical properties of matter. For observational consistency it is assumed that λ is of order (63). A dynamical derivation is possible if one goes into the very early inflationary epochs. However, the values of λ deduced from those calculations are embarrassingly large, being $10^{108} - 10^{120}$ times the value given by (63). The problem then is, how to reduce λ from such high values to the presently acceptable range (Weinberg 1989). By contrast, the present derivation leads to the acceptable range of values with very few theoretical assumptions.

The theory developed in this paper differs from big-bang cosmology in what we believe to be an important aspect, that the gravitational equations are scale invariant. The gravitational equations including both the creation terms and the cosmological constant then reduce in the constant mass frame to

$$R_{ik} - \frac{1}{2}g_{ik}R + \lambda g_{ik} = -8\pi G\left[T_{ik} - \frac{2}{3}(c_i c_k - g_{ik}\frac{1}{4}c_l c^l)\right]. \tag{64}$$

The immediate successes of the theory are :

(i) The circumstance that G determined by (34) is necessarily positive requires gravitation to act as an attractive force. Aggregates of matter must tend to pull together. This is unlike general relativity where gravitation can as well be centrifugal, with aggregates of matter blowing always apart, as follows if G in the action (3) of general relativity is chosen to be negative.

(ii) In the cosmological case with homogeneity and isotropy the pressure contributed by the c-field term in the gravitational equations is negative, explaining the expansion of the universe.

(iii) Also in the cosmological case, the energy contribution of the c-field is negative, which ensures that when the creation conditions (9) are satisfied the creation process tends to cascade with explosive consequences.

(iv) The magnitude of the constant λ is shown to be of the order needed for cosmology. Unlike big-bang cosmology this is a deduction not an assumption.

These ideas therefore generate hopes for a more comprehensive framework for relating the property of inertia of matter and the phenomena of matter creation to cosmology. It is not claimed that what is outlined here is the final product; rather it should be looked upon as a preliminary attempt. The successes claimed above have to be followed up by a quantum version of the Machian theory and the empirical values of the parameters of the quasi-steady state model have to be related to the fundamental constants of the theory as well as to cosmological boundary conditions. This is the direction in which our future theoretical work will go.

3. Observational Facts and Consequences

Earlier we showed that the approximate solution for the scale factor $S(t)$ is given by equation

$$S(t) = \exp\left(\frac{t}{P}\right)\left\{1 + \alpha\cos\frac{2\pi t}{Q}\right\}.$$
$$\text{where } P \gg Q \tag{22}$$

We have chosen values of Q and P as follows.

$$Q = 4 \times 10^{10} yr,\ P = 20Q \tag{23}$$

In (22) α, P, Q are constants determining the model, with $P \gg Q$ a consequence of creation being slow. We now examine the astrophysical consequences of the scale factor $S(t)$ being given by (22).

We are thus concerned with an oscillatory model in which some matter creation occurs, especially near the minimum in each cycle, as was already visualized in HBN 1993. At each oscillation the universe experiences an expansive push. To give a framework for discussion, we suppose creation occurs so that the ratio S_1/S_2 stays fixed, as (22) requires it to do, with S_1

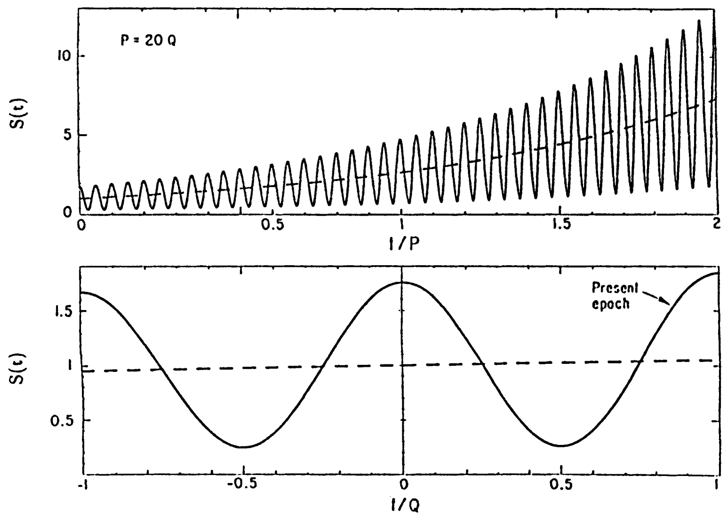

Figure 1. S(t) plotted against t/P (upper panel) and against t/Q (lower panel).

and S_2 both increasing as the slow exponential factor $exp\ t/P$. Thus the timescale for the universe to expand irreversibly by e is $P \gg Q$, which is to say in each exponentiation there are many oscillations. The situation is analogous to the classical steady-state model but with each exponentiation of the scale factor broken into many oscillations.

In Fig 1 we show $S(t)$ plotted against t/P and against t/Q for the assumed value of $P = 20Q$. We also put $\alpha = 0.75$. The time in Fig 1 is measured in units of Q. In order to relate this model to the current state of the observed universe we also need to assign a value for t_o, the present epoch, in relation to the phase of the oscillatory cycles. We choose $t_o = 0.85$ being 85 per cent of the way through the current cycle, cycles being reckoned maximum to maximum.

The parameter Q is related to the observed values of the Hubble constant H_o and the deceleration term q_o respectively. For $P \gg Q$, the effect on H_o, q_o of the overall expansion will hardly be noticed. The time dependent quantities H, q defined as

$$H = \frac{\dot{S}}{S}, \; q = -\frac{S\ddot{S}}{\dot{S}^2} \tag{65}$$

have the following properties. Starting from the minimum phase of an oscillation, H begins at zero, rises to a maximum and then falls back to zero at maximum phase, while q starts sharply negative and grows to zero, and then

TABLE 1.

Q(years)	$H_o(km\ sec^{-1}\ Mpc^{-1})$
30×10^9	86.2
40×10^9	64.7
50×10^9	51.7

goes to markedly positive values as maximum phase is approached. The observed value of H, H_o, lies between $\sim 50 km\ sec^{-1}\ Mpc^{-1}$ (cf Sandage 1993) and $\sim 80 km\ sec^{-1}\ Mpc^{-1}$ (cf Tully 1993) while Kristian, Sandage and Westphal (1978) gave $q_o \simeq 1.5$ but with considerable uncertainty. These values are generally indicative of a phase in the current oscillation approaching maximum. Without knowing the precise present-day phase, the time that has elapsed since the last minimum is somewhat uncertain but is probably close to $\frac{1}{2}H_o^{-1}$.

This leads us to the following numerical values relating Q to H_o.

In HBN 1994a we chose the value $Q = 40 \times 10^9$ years with $H_o = 64.7 km\ sec^{-1} Mpc^{-1}$ as a compromise between the high and low values of H_o which are a subject of continuous debate. The parameter q_o is determined from (22) to be 1.725, again close to the value given above. With the choices of α and t_o the maximum redshift of objects in the present cycle is $z = 4.86$. This is not a limit, however, since the corresponding redshift from the previous cycle will be $z = 5.166$ and so on step by step. However the Hubble diagram for this model shows that the objects in each cycle will be fainter than those in the most recent cycle by about 3 magnitudes. Thus by specifying α, t_o and the parameter Q based on the observed value of H_o, we obtain reasonable values for q_o, the maximum redshift in this cycle, and a Hubble diagram.

In this model we also have to explain two other properties of the observed universe which have previously been thought to provide strong evidence for the so-called standard model.

3.1. THE COSMIC MICROWAVE BACKGROUND

It has been known for many years that the energy density of the microwave background is almost exactly equal to the energy released in the conversion of hydrogen to helium in the visible baryonic matter in the universe (cf Hoyle 1968). This density is $\rho \simeq 3 \times 10^{-31} gm\ cm^{-3}$ and we suppose that about $7.5 \times 10^{-32} gm\ cm^{-3}$ is H_e. Thus the energy released in the production of this H_e through the conversion $H \rightarrow He$ is $4.5 \times 10^{-13} erg\ cm^{-3}$, which

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if thermalized gives a radiation field of $2.78K$.

In the standard Big-Bang cosmology this agreement with the observed value is considered to be purely fortuitous, but within the framework of QSSC it is a clear indication that the microwave background was generated ultimately by the burning of hydrogen into helium in stars, through many creation cycles each of length Q. The optical and ultraviolet light must have progressively been degraded and scattered by dust, much of it in the form of iron needles so that it now forms a smooth black body form, as discussed at length in HBN 1994, where we predict a temperature of $2.68K$, very close to the observed temperature of $2.735 \pm 0.06K$ (Mather et al. 1990).

How many cycles are required, i.e. what is the value of P/Q? We have shown in HBN 1994 that the ratio P/Q can be obtained from the observed $\log N - \log S$ curves for radio sources. This is because radio sources from earlier cycles are contained in the counts. The reason for this is that while there will be optical obscuration near oscillatory minima and optical sources from earlier cycles will not be easily detectable, this will not apply at long enough radio wavelengths. For a simple model in which it is assumed that radio sources appear at a uniform rate per unit proper volume, we find that $P/Q \simeq 20$. With this value we not only can understand the origin of the microwave background but also the shape of the $\log N - \log S$ radio observations to very faint levels. Thus we have shown that in this model we have two timescales $Q = 40 \times 10^9$ years and $P = 8 \times 10^{11}$ years.

3.2. PRODUCTION OF THE LIGHT ISOTOPES

In the standard model the production of the light elements is attributed to nuclear reactions early in the explosion. In HBN 1993 (Section 6 and Appendix) in Hoyle (1992), and most recently in Hoyle, Burbidge and Narlikar (1995),(HBN 1995), a detailed analysis has been given of a similar process in QSSC in which the light elements are synthesized in a creation process starting with a Planck fireball. In view of the fact that remarkable results are obtained using this approach, we describe it again in some detail here.

Because the early stages in the development of a Planck fireball belong to the realm of unknown physics, it is necessary to begin with a specification of initial conditions. Fermions of familiar types are necessarily excluded by degeneracy conditions at early stages when the fireball dimension is only $\sim 10^{-33}cm$. Indeed, fermions of familiar types cannot appear until the interparticle spacing within the expanding fireball has increased to $\sim 10^{-13}cm$.

We take the view in specifying the model to be investigated that energy considerations discriminate against charmed, bottom and top quarks. We also take the view that degeneracy considerations, together with the need

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TABLE 2. Densities and Temperatures at $1 < r < 4$ in the expansion of a Planck Fireball

r	1	1.25	1.5	1.75	2	2.5	3	3.5	4
$\log N$	36.08	35.79	35.55	35.35	35.18	34.89	34.65	34.45	34.27
T_9	0	13.9	19.3	21.2	21.7	20.8	19.3	17.7	16.3

for electrical neutrality, prevent the strange quark from being discriminated against. When the up, down and strange quarks combine to baryons, equal numbers of N, P, Λ $\Sigma^{\pm}$, Σ^0, Ξ^0, Ξ^- are thus formed, with only a negligible amount of Ω^-. Because of the long lifetimes, $\sim 10^{-10}$ seconds, of Λ, $\Sigma^{\pm}$, Ξ^0, Ξ^-, the strange quark survives the effective stages in the expansion of the fireballs, although Σ^0 goes to Λ plus a γ-ray at a stage proceeding the synthesis of the light elements. Finally, we consider that baryons containing the strange quark do not form stable nuclei. Ultimately they decay into N and P, but only after the particle density has fallen so far that the production of light elements has stopped. With N going on a much longer time scale (10 minutes) into P, six of the baryons of the octet go at last into hydrogen. Thus we see immediately that the fraction by mass of helium, Y, to emerge from Planck fireballs is given by

$$Y = 0.25(1-y), \tag{66}$$

where $1-y$ is the fraction of the original N and P to go to ^4He. Anticipating that y will be shown in the next section to be ~ 0.085, equation (66) gives $Y = 0.229$, somewhat lower than the value of ~ 0.237 obtained previously (Hoyle, 1992).

The numerical values used in the detailed calculations of later sections are given in Table 2.

Here N is the number per cm^3 of each baryon type, the values in the table being such that N declines with increasing r as r^{-3}. The unit of r depends on a specification of the total number of baryons in the fireball. Thus for a total of 5.10^{18} the unit of r is $5.10^{-7} cm$. However, since this total is uncertain, because the Planck mass, usually given as $(3\hbar c/4\pi G)^{1/2}$, is theoretically uncertain to within factors such as 4π, we prefer to leave the unit of r unspecified – we shall not need it in the calculations. Suffice it that there will always be a unit for r such that N has the values in the table.

Taking the expansion of the fireball to occur at a uniform speed v, the time t of the expansion to radius r is proportional to r, $t \propto r$. In specifying the model we take the factor of proportionality here to be 10^{-16} seconds. With the unit of r chosen as $5.10^{-7} cm$ this requires $v = 5.10^9 cm\ s^{-1}$, a rather low speed. But for a Planck mass increased by 4π above $(3\hbar c/4\pi G)^{1/2}$ the expansion speed is raised by $(4\pi)^{1/3}$ to $1.16 \times 10^{10} cm\ s^{-1}$. Thus

$$t = 10^{-16} r \text{ seconds,} \tag{67}$$

thereby relating t to N and T_9 through the values in Table 2. The numerical coefficient of equation (2) can be regarded as a parameter of the theory, but it is not a parameter that can be varied by more than a small factor, unlike the parameter η in Big-Bang nucleosynthesis which could be varied by many orders of magnitude for all one knows from the theory.

The temperature values in Table 2 are calculated from a heating source which comes into play at $r = 1$, i.e. at $t = 10^{-16} s$. The heating source is from the decay of π^0 mesons with a mean life of $8.4 \times 10^{-17} s$. The temperature values in Table 2 correspond to a π^0 meson concentration of $2/3N\ cm^{-3}$, which is to say one π meson to each neutron and each proton, with π^0, $\pi^{\pm}$ in equal numbers.

The decay of a π^0 meson into two 75 Mev γ-rays does not immediately deposit energy into the temperature T_9 of the heavy particles. It does not even lead to more than a limited production of $e^{\pm}$ pairs, because at these densities this is prevented by electron degeneracy. Thus the energy of π^0 decay is at first stored in the form of relativistic particles, quanta and some $e^{\pm}$, the latter being adequate, however, to prevent the γ-rays from escaping out of the fireball.

As the fireball expands, confined relativistic particles lose energy proportional to $1/r$, the energy loss going to the heavy particles, for each type of which there is a conservation equation of the form $dQ = dE + PdV$, viz

$$-\alpha d(1/r) \quad = \quad 3/2 k dT + 3kT dV/V, \tag{68}$$

an equation that integrates to give

$$T_9 \quad = \quad \frac{2}{3}\frac{\alpha}{k}\ \frac{r-1}{r^2}, \tag{69}$$

the constant of integration being chosen to given $T_9 = 0$ at $r = 1$. The constant α in is easily determined from the energy yield of the π^0 mesons. Sharing the energy communicated to the heavy particles equally among all of them, leads to the values of T_9 in Table 2.

The energy is considered to have all gone to the heavy particles by the stage of the expansion when r reaches 4, after which T_9 declines as r^{-2}, i.e. adiabaticlly, the heavy particles being non-relativistic in their thermal motions. Thus for $r > 4$ we have

$$T_9 = 16.3. \left(\frac{4}{r}\right)^2 = \frac{260.8}{r^2}, \tag{70}$$

$$t = 10^{-16} r = \frac{1.62 \times 10^{-15}}{T_9^{1/2}} \text{seconds}, \tag{71}$$

while the particle densities decline as r^{-3}.

(a) The Abundance of ^4He

It will be shown that neutrons and protons are in statistical equilibrium with ^4He up to $r = 3$ in Table 2, but not for $r > 3$. Defining a parameter ζ by

$$\log \zeta = \log N - 34.07 - \frac{3}{2} \log T_9 \tag{72}$$

it was shown by Hoyle (1992) that the fraction y of neutrons and protons remaining free at temperature T_9 and particle density N for each nucleon type is given in statistical equilibrium by

$$\log \frac{1-y}{y^4} = 0.90 + 3 \log \zeta + \frac{142.6}{T_9}, \tag{73}$$

the values of T_9 and N in Table 2 at $r = 3$ giving $y = 0.085$, leading to the value $Y = 0.229$ given above. A similar calculation at $r = 2.5$ yields $y = 0.083$, much the same as at $r = 3$. For $r < 2.5$ the values of y fall away to ~ 0.06. Thus in moving to the right in the table the values of y increase towards $r = 3$, where the falling value of T_9 eventually freezes the equilibrium.

The condition for freezing is that the break-up of ^4He by $^4He(2N, T)T$, followed by the break-up of T and D into neutrons and protons should just be capable of supplying the densities of P and N, $n(P) = n(N) \simeq 5.10^{33} cm^{-3}$ for the range of r from 2.5 to 3 and $y \simeq 0.085$. The time available for this break-up of ^4He is that for r to increase from ~ 2.5 to ~ 3, i.e. 5.10^{-17} seconds. In this time the break-up of $n(A) \simeq 2.9 \times 10^{34} cm^{-3}$, Using the reaction rates of Fowler, Coughlan and Zimmerman (1975) we verify that.

$$\frac{1.67 \times 10^9}{T_9} . \frac{3.28 \times 10^{-10}}{T_9^{2/3}} \exp - \frac{4.872}{T_9^{1/3}} . \exp - \frac{131.51}{T_9}$$
$$(1 + 0.086 T_9^{1/3} - 0.455 T_9^{2/3} - 0.271 T_9 + 0.108 T_9^{4/3} + 0.225 T_9^{5/3}$$
$$\left(\frac{n(N)}{6.022 \times 10^{23}}\right)^2 n(A) . \, 5 \times 10^{-17} \qquad (74)$$

Here we put $T_9 \simeq 20$ for the range of r from 2.5 to 3, and putting $n(N) = 5.10^{33} cm^{-3}$, $n\ (A) = 2.9 \times 10^{34} cm^{-2}$, the value of (74) is $2.85 \times 10^{33} cm^{-3}$, This is close enough to the required value of 2.5×10^{33}.

This is already an astonishing result. That so complicated an expression as (74) should combine so exactly to produce such an outcome is not a consequence of the parametric choice of the model. The freedom of choice of the numerical coefficient in (74) is entirely dwarfed by the factors 10^{34}, 10^{33}, 10^9, 10^{-10}, 10^{-17} in (74), while even some variation in the parameter α in (4), as it affects the value of the factor $\exp -131.51/T_9 \approx 2.5 \times 10^{-3}$, is also dwarfed by the much larger powers in (74). The most license that can be permitted to a critic would be to accept the above result as model-dependent to the extent that it already consumes essentially all the available degrees of freedom of the model, leaving all further results to be judged as effectively parameter independent.

(b) The Abundances of D and ^3He

We have omitted the analysis which leads to the values

$$D/H =^3 He/H \simeq 5 \times 10^{-5} \qquad (75)$$

which are given in the summary Table 3.

(c) The Abundance of ^7Li

Writing $n(P)$, $n(A)$ for the densities of protons and alpha particles we have

$$n(P) \quad = \quad 1.58 \times 10^{33} \left(\frac{T_9}{16.3}\right)^{3/2} cm^{-3}, n(A) = 8.5 \times 10^{33} \left(\frac{T_9}{16.3}\right)^{3/2} (76)$$

The ratio of the abundance of ^7Li to ^8Be established in statistical equilibrium at temperature T_9 is given by

$$\log \frac{^7Li}{^8Be} = \frac{3}{2}\log\frac{7}{8} + \log 4 - \log n(P) + 34.07$$
$$+\frac{3}{2}\log T_9 - \frac{5.04}{T_9} \times 17.35,$$
$$= 3.20 - \frac{87.44}{T_9}, \quad (77)$$

with

$$\log \frac{^8Be}{^4He} = \frac{3}{2}\log 2 + \log n(A) - 34.67 - \frac{3}{2}\log T_9 = -2.11 \quad (78)$$

also given by statistical equilibrium.

The abundance of ^7Li established at T_9 according to (23) will, however, be subject to attenuation as the temperature declines from T_9, according to an attenuation factor

$$A\int_0^{T_9} \exp -\frac{30.443}{T_9} dt, \quad (79)$$

with

$$dt = \frac{8.1 \times 10^{-16}}{T_9^{3/2}} dT_9 \quad (80)$$

as before and A a numerical coefficient obtained from the reaction rate for $^7Li(P,A)^4He$ given by Fowler et al. (1975), viz

$$A = 1.7 \times \frac{1.05 \times 10^{10}}{T^{3/2}} . \frac{2.40 \times 10^{31}}{6.022 \times 10^{23}} T^{3/2} = 7.25 \times 10^{17} s^{-1} \quad (81)$$

The factor 1.7 here arises from an estimate of the combined effect of various terms adding to the rate of $^7Li(P,A)^4He$, the rest of A being the main term. Evaluating (79) leads to

$$\sim 19.3 T_9^{1/2} \exp - \frac{30.443}{T_9} \tag{82}$$

as the attenuation factor to be applied to the abundance of ^7Li. With $\log \frac{^4He}{H} = -1.08$ we thus have

$$\begin{aligned} \log \frac{^7Li}{H} &= 3.20 - 2.11 - 1.08 \\ &-19.3 \times 0.4343 T_9^{1/2} \exp - \frac{30.443}{T_9} \\ &= 0.01 - 8.38 T_9^{1/2} \exp - \frac{30.443}{T_9} \end{aligned} \tag{83}$$

which has a maximum of -9.60 at $T_9 \simeq 12$. Thus the surviving lithium abundance is

$$\frac{^7Li}{H} \simeq 2.50 \times 10^{-10}, \tag{84}$$

a result in good agreement with the observational requirement, again calculated from highly complicated formulae, again without any model adjustment.

(d) The Abundance of ^{11}B

A similar calculation for ^{11}B leads to $^{11}B/H \simeq 10^{-18}$, below the observational detection limit. This is significantly lower than the value calculated by Hoyle (1992) who used an attenuation factor that was not sufficient. From an observational point of view the model therefore predicts that there is effectively no 'plateau' under boron. Such boron as exists is required to come from cosmic-ray spallation on ^{12}C, ^{16}O.

(e) The Abundance of ^9Be

As noted in Hoyle (1992), the nucleus of ^9Be is exceptionally fragile, leading to a particularly low freezing temperature. Statistical equilibrium at higher temperatures establishes

$$\log \frac{^9Be}{H} = \frac{3}{2} \log \frac{9}{8} + \log \frac{4}{3} - 0.15 + \log D/H$$

$$+\log\frac{^8Be}{^4He}+\log\frac{^4He}{n(P)}-\frac{3.28}{T_9} \qquad (85)$$

with respect to the reaction $^9Be(P,D)2^4He$. Using $\log D/H=-4.30$ already calculated, $\log\, ^8Be/^4He=-2.11$, $\log^4 He/n(P)=0.73$, gives $-5.63-3.28/T_9$ for the right hand side of (?). Because $^9Be(P,A)^6Li$ contributes equally with $^9Be(P,D)2^4He$ to the destruction of ^9Be, whereas at $T_9\simeq 1$ it contributes essentially nothing to the production of ^9Be, the equilibrium concentration of ^9Be is lowered by a further factor 2, so that

$$\log\frac{^9Be}{H} = -5.93-\frac{3.28}{T_9}. \qquad (86)$$

Freezing of the equilibrium condition at $T_9=0.50$ for ^9Be would thus give

$$\log\frac{^9Be}{H} = -12.5 \qquad (87)$$

in satisfactory agreement with the apparent observed plateau under ^9Be (Boesgaard, 1994).

The estimated freezing temperature according to the model can be obtained by requiring that the product of the expansion time scale, $1.62\times 10^{-15}/T_9^{1/2}$ seconds at temperature T_9 and the sum of the reaction rates terms for $^9Be(P,D)2\,^4He$ and of those for $^9Be(P,A)^6Li$ be unity, viz

$$2.\frac{1.03\times10^9}{T_9}\cdot\frac{2.40\times10^{31}}{6.033\times10^{23}}.\frac{1.62\times10^{-15}}{T_9^{1/2}}T_9^{3/2}.\exp-\frac{3.046}{T_9} = 1. \qquad (88)$$

The factor 2 on the left of this formula comes from the circumstance that at the values of T_9 in question the highly complicated non-resonant contribution given by FCZ about doubles the resonant reaction rates. Equation (88) determines a freezing temperature $T_9=0.623$, reasonably close to the required value of 0.5.

(f) The Abundances of ^{12}C and ^{16}O

The reaction rate on ^9Be from $^9Be(A,N)^{12}C$ as given by FCZ is

$$\sim\frac{2.40\times10^8}{T_9^{3/2}}\frac{n(A)}{6.023\times10^{23}}.\exp-\frac{12.732}{T_9}. \qquad (89)$$

Using (74) for $n(A)$ and putting $T_9 \simeq 10$, at which temperature most of the production of ^{12}C takes place, gives 1.44×10^{16} for (89). Multiplying by the time-scale $1.62 \times 10^{-15}/T_9^{1/2}$ then gives ~ 7.4, implying that an abundance $^9Be/H \simeq 5.5 \times 10^{-7}$ given by (20) is converted 7.4 times over to ^{12}C, leading to

$$\frac{^{12}C}{H} \simeq 4.1 \times 10^{-6} \tag{90}$$

The value of $^{16}O/H$ is of a similar order.

(g) The External Medium

All of the above followed from just the N and P members of the baryon octet. The other six baryons are considered not to form stable nuclei. They decay in $\sim 10^{-10}$ seconds, by which time a Planck fireball has effectively expanded into its surroundings, which according to the QSSC model (HBN 1993,1994a,b) is necessarily a strong gravitational field in which the decay products of Λ, $\Sigma^{\pm}$, Ξ^0 and Ξ^- may be expected rapidly to lose energy. The Ξ^0 baryon decays to Λ and π^0 in a mean life of $3.0 \times 10^{-10}s$, Σ^+ which decays in a mean life of $8.0 \times 10^{-11}s$, gives a π^0 meson in about a half of the cases, so that together with Λ, which decays in a mean life of $2.5 \times 10^{-10}s$, there is a late production of about $2.5\pi^0$ per baryon octet, yielding ~ 5 late γ-rays per octet, typically with energies ~ 100 Mev. It is these γ-rays and their products that are expected to be subjected to energy loss in strong gravitational fields.

The production of Planck particles near large masses of the order of galactic clusters occurs typically in an environmental density $\sim 10^{-16} g\ cm^{-3}$ at which density γ-rays of 100 Mev have path lengths of $\sim 10^{18} cm$, ample for considerable redshifting effects to occur, when quanta and particles in the $1 - 100$ kev range would arise. Although such particles and quanta are readily shielded against, it is an interesting speculation that pathways into the external universe may be briefly opened and that the mysterious γ-ray bursts arise in such situations.

(h) Summary of Abundances and Conclusions

The calculations more accurate than those described earlier in Hoyle (1992) and in HBN (1993). They lead to the abundances and results shown in the following table.

To obtain a ratio $^9Be/H \simeq 3.10^{-13}$ requires a freezing temperature $T_9 \simeq 0.5$ which is close but not equal to the calculated freezing temperature $T_9 \simeq 0.62$.

TABLE 3. Summary of Results

$^4He/H$	$=$	$Y \simeq 0.229$
$D/H \simeq ^3He/H$	$\simeq$	5.10^{-5}
$^7Li/H$	$=$	2.5×10^{-10}
$^{11}B/H$		very small
$^{12}C/H$	$\simeq$	4.1×10^{-6}

We conclude that a certain model of the decay of Planck particles leads to interesting values for the abundances of the light elements. The work is deductive, and in this sense the model used is not subject to negotiation, any more than the axioms on which a mathematical theorem is proved are subject to negotiation. Or any more than supporters of Big-Bang nucleosynthesis regard the choice of their parameter η as a matter of negotiation. Thus the only basis for judging the situation is to assess how good, or bad, are the results. Our assessment is the following:

(i) Our result for $^4He/H$ is very good.

(ii) The ratios D/H, $^3He/H$ are too high by factor ~ 2. A more detailed calculation might well lower $^3He/H$ to its observational value. But at the expense of a further increase in D/H, necessitating an epicycle for the theory in which the observed $D/H \simeq 1.5 \times 10^{-5}$ is due to environmental effects.

(iii) The ratio $^7Li/H$ is very good.

(iv) The prediction of essentially no 'floor' under ^{11}B is subject to test. The 'floor' under ^9Be requires a freezing temperature $T_9 \simeq 0.50$, whereas the calculated freezing temperature was $T_9 \simeq 0.62$. Considering the very complicated expressions of FCZ, especially that involved in a cut-off procedure for non-resonant contributions, this correspondence is adequately close.

Finally we may ask how this situation for the synthesis of light elements from Planck particles compares with the situation in Big-Bang nucleosynthesis. In that case

(a) The classic choice $\eta = 3.10^{-10}$ for the baryon to photon ratio is good for $^3He/H$ but is too low for $^7Li/H$ and too high for Y and D/H.

(b) While reducing η brings Y and $^7Li/H$ into good agreement with observation the value of D/H becomes so large that the theory requires an astration epicycle to save itself.

(c) Raising η to $\sim 6.10^{-10}$ gives good results for D/H, $^3He/H$ and $^7Li/H$ but the resulting value $Y = 0.25$ is too high, and hardly savable by any epicycle or combination of epicycles.

(d) The theory predicts no plateau under ^9Be, which seems wrong. A recourse to inhomogeneous cosmological models would be to make the theory wildly epicyclic.

(e) Big-bang nucleosynthesis, but not the present model, predicts a present-day average baryon density in the universe much below the cosmological closure value, either forcing a change to a so-called open model (when galaxy formation is made difficult or impossible) or leading to the proposal that most of the material in the universe must be dark and non-baryonic. It is this argument that has led to the proposal that non-baryonic matter dominates the universe. None has so far been found.

3.3. THE VALUE OF $\rho_\circ$

To determine $\rho_\circ$ (the mean density at this epoch) we have no simple relation like $\rho_\circ = 3H_\circ^2/8\pi G$ in the closure model of Friedmann cosmology. However from the analysis in Part I we have that

$$\frac{\dot{S}^2}{S^2} = \frac{8\pi G\rho_\circ}{3}\left(\frac{S_\circ}{S}\right)^3 + \frac{1}{3}\lambda\,. \qquad (91)$$

Also neglecting the slight variation of $exp\ t/P$ over the current half-cycle,

$$\frac{S_\circ}{S} = \frac{1 + 0.75\cos 1.7\pi}{1 + 0.75\cos 2\pi t} \qquad (92)$$

in which $\alpha = 0.75$ and $t_\circ = 0.85$ are used. Applying (91) and (92) at $t = 1$, the next maximum when $\dot{S} = 0$, the value of $\rho_\circ$ is related to λ by

$$\lambda = -0.558.\, 8\pi G\rho_\circ\,. \qquad (93)$$

Substituting λ given by (93) in (91) now gives

$$\frac{\dot{S}^2}{S^2} = \frac{8\pi G}{3}\rho_\circ\left[\left(\frac{S_\circ}{S}\right)^3 - 0.558\right]. \qquad (94)$$

Since $\dot{S}^2/S^2$ at $t = t_\circ$ is $H_\circ^2$ we therefore get

$$0.442\rho_\circ = \frac{3H_\circ^2}{8\pi G}, \qquad (95)$$

the coefficient 0.442 being appropriate only for the present moment $t_\circ = 0.85$. Putting $H_\circ = 64.7 km\ sec^{-1} mpc^{-1}$, $(Q = 40 \times 10^9)$, determines the present-day average cosmological density as

$$\rho_\circ \quad = \quad 1.79 \times 10^{-29} g\ cm^{-3}\ . \tag{96}$$

Thus, since the density of visible matter is $\approx 3 \times 10^{-31} gm\ cm^{-3}$ in this theory as in the standard big bang the bulk of the matter is dark. However the major difference between the QSSC and the big bang is that the dark matter in our theory is made up of baryons. This points up the fact that it is only in the big bang scenario that there is any reason at all to suppose that dark matter is non-baryonic.

3.4. THE COMPOSITION OF THE DARK MATTER

Since a typical galaxy goes through many cycles Q in our model, stellar evolution is not limited by the value of $2/3H_0 \simeq 15 \times 10^9$ years as it is in the big bang.

In our Galaxy we then will expect to have the following stellar components:

1. The known stellar population with stars with ages ranging from $\sim 10^7$ years for those most recently formed to $\sim 15 \times 10^9$ years for the oldest globular clusters.
2. Stars which were formed from matter created in a previous cycles, and are now $\sim 50 \times 10^9$ years old, and remnants from even earlier cycles.
3. Failed stars – the so-called MACHOS.

The distribution of all of this baryonic matter in disk and halo is determined by the formation process and its dynamical behavior. Creation events on a scale sufficient to augment the masses of galaxies on a significant scale will produce violent gravitational disturbances, mostly disrupting and expelling into a halo all previously existing stars. Thus the present day galactic disk (Population I) is only the most recent of a number of major star-forming episodes. Previous episodes occurring at ages $(15 + 50P) \times 10^9$ years, P–1,2,3,..., are now in the halo. Indeed they form the halo, and the process just described gives a picture of how the halos of galaxies are built up. All galaxies have a largest value of P, corresponding to the epoch of their formation, the larger P_{max} corresponding to massive ellipticals and the smaller P_{max} corresponding to 'late' type spirals and irregulars. We estimate $P_{max} \sim 10$ for a galaxy such as our own and $P_{max} \sim 20$ for a typical massive elliptical.

The main-sequence of a stellar population of age 15×10^9 years is burnt out down to masses of solar order and is usually taken to have mass to light ratio of ~ 3, while a population of age $(15+50P)\times 10^9$ years will have its main-sequence burnt-out down to a mass $\sim M_\odot(\frac{1+10P}{3})^{\frac{1}{4}}$ and will have a mass to light ratio $\sim 3\times(1+10\ P/3)$ which for $P=10$ have the values $0.41M_\odot$ and ~ 100 respectively.

The stellar mass function determined in the Galaxy from the solar neighborhood has a 'missing mass' factor of about 2. While this is uncertain, it can be explained by the addition of so-called brown dwarfs to the mass function. If the stellar mass function is taken to be everywhere and always the same, high mass-to-light ratios for halo populations and for elliptical galaxies cannot be explained wholly through a brown dwarf component. But much or most of such high mass-to-light ratios can be explained if such populations have ages of $(15+50P)\times 10^9$ years with $P\sim 10$ or more.

It is a prediction of our theory that the main-sequence of halo stars should thicken for red dwarfs with masses $\sim 0.5M_\odot$. Attempts to confirm or deny this prediction would need to be confined, because of the low luminosities of such dwarfs, to the solar neighborhood, say to with $\sim 50pc$ of the sun. Being halo stars they would pass through the solar neighborhood with high velocities, $\sim 200km\ s^{-1}$. An accurate color-magnitude for the high velocity stars of the solar neighborhood would therefore be of great cosmogonic interest. To which should be added the need for accurate theoretical studies of the evolution of very old stars with masses $\sim 0.5M_\odot$, especially with regard to the possibility of the long-term mixing of the products of nucleosynthesis.

3.5. FAINT GALAXIES

The apparent luminosity of a galaxy of radial co-ordinate r and intrinsic luminosity L observed at a redshift z is given by

$$\sim \frac{L}{4\pi r^2}\frac{1}{(1+z)^2}\frac{1}{S^2(t_\circ)} \tag{97}$$

Putting $rS(t_\circ)\simeq cH_\circ^{-1}\simeq 2\times 10^{28}cm, z=z_1=5$, a galaxy of absolute magnitude –21 would thus be observed with an apparent bolometric magnitude of about +27. Since this is within the range of observation it follows that galaxies with still larger values of r should also be observable, such galaxies having an emission time t that occurred in the previous universal oscillation. For those where emission occurred at the last oscillatory maximum there would be a blueshift with $S(t)>S(t_\circ)$. If we suppose that at the present we are not far from maximum phase, the blueshift will be comparatively small, and the second factor in (97) will be greater than unity but not greatly so. With $rS(t_\circ)\simeq 5\times 10^{28}cm$ in such a case, the apparent bolometric magnitude of a galaxy of absolute magnitude –21 would be

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somewhat fainter than +26 if absorption of a magnitude at the last oscillatory minimum is included. The theory thus predicts that a multitude of blue galaxies should be observed at about this brightness level, some indeed with spectrum lines that are blueshifted rather than redshifted. The blueness is not an intrinsic property of the galaxies themselves but arises from the oscillatory character of the scale function $S(t)$ together with there being many oscillations occurring in the characteristic expansion time $P \simeq 10^{12}$ years of the universe, and with there being little change of the universe from one oscillation to the next.

Thus the prediction is that faint blue galaxies will appear in profusion at faint magnitude levels. Other explanations of this *observed* phenomenon have been proposed (cf Koo and Kron 1992) but here we have an explanation which comes naturally out of the cosmological model.

Also in this model the universe changes much more slowly than had hitherto been supposed since the appropriate timescale is determined by P. Some 5-10 exponentiations of the scale factor $S(t)$ are required to expand an initially local situation with a dimension of a few megaparsecs to dimensions $\sim 3000 Mpc$. In terms of the galaxies we observe, the average age is $P/3 \sim 3 \times 10^{11}$ years, while the ages of the oldest objects at the limit of observation are $5 - 10P \approx 10^{13}$ years. Clusters of ellipticals like the Coma cluster may well have ages intermediate between these values, i.e. $\sim 2.10^{12}$ years. On this basis we would expect that a large part of the mass will be in the form of evolved stars, not only white dwarfs, neutron stars or black holes but also dead stars with $M > 0.5 M_{\odot}$, including brown dwarfs. Some part of this may be in the form of completely evolved galaxies.

3.6. THE COSMOGONY ASSOCIATED WITH THE QSSC

The long time scale associated with this cosmology means that galaxies can form at all epochs and that different components in the same galaxy will have very different ages. Thus stellar components with ages in the range $10^7 - 10^{12}$ years will be present.

Since the mass creation events will be at maximum intensity and frequency in the minima of the cycles and decrease in intensity and numbers at the peaks of the cycles young objects will be comparatively rare among the galaxies at comparatively small redshifts.

Observational evidence of a range of ages is present. The evidence is of several different kinds.

(a) A class of faint galaxies often called H II galaxies which were originally investigated by Sargent and Searle (1970) and Kunth and Sargent (1983) are found among the faint galaxies.. Their spectra have very weak continua and line emission characteristics of H II regions together with O

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and B stars. All of the strong features in the spectra arise from stars with ages $\leq 10^8$ yr. There is some ambiguity about the continua. In one case, Sandage (1963) argued that the colors of the continuum in NGC 2444-2445 were similar to those in the Large Magellanic Cloud and that there could be an underlying old system. However, in general there is no strong evidence that any old stellar population is present.

b) Highly luminous *IRAS* galaxies: Far-infrared observations ($\sim 25 - 200\mu m$) made originally in a few cases from the ground or high-flying aircraft, and then much more extensively from *IRAS*, have shown that there is a large population of spiral and irregular galaxies which emit powerfully in the far-infrared (Soifer, Houck, & Neugebauer 1988). It is generally believed that the far-infrared radiation is thermal emission from dust heated by main-sequence stars, and that the powerful sources are regions in which star formation is dominant – so-called starburst regions. The most extreme examples are the high-luminosity *IRAS* galaxies with luminosities in excess of $10^{12} L_{\odot}$, practically all in the far-infrared. Nearby examples of such galaxies are Arp 220 and NGC 6240. Such systems are all very irregular. While it has often been argued that such objects are a result of mergers between previously separate galaxies (cf. monograph edited by Wielen 1989), one of us (Burbidge 1986; Burbidge & Hewitt 1988) has made the case that these objects are genuinely young ($\ll 10^9$ years old) systems made up of successive generations starting with fairly massive stars ($20 M_{\odot} \leq M \leq 50 M_{\odot}$) which have themselves in the early generations made the dust which is now being heated by ultraviolet radiation from later generations. These galaxies fulfill all of the tests for young systems. There are no stars detected in them older than A-type systems, they contain huge masses of molecules and we predict that the dust in them will not be of typical galactic form, since it will have condensed from metallic oxides and other compositions typical of material ejected from massive stars. According to our proposal made here, galaxies of the types found in categories *a* and *b* have arisen from recent mass-creation events. Another unexpected observational discovery has been the finding of very young stellar systems in old galaxies. There are many examples. We mention here two recent discoveries.

c) In our own Galaxy, Krabbe et al. (1991) have shown that within 0.5 pc of the center there is a cluster of young massive stars with random motions of $\sim 200 km\ s^{-1}$. These have been found by high-resolution imaging of a $2.06\mu m$ recombination line of He I. The stars are broad emission line objects which must have ages less than 10^6 yr. The radio source Sgr A lies in this cluster. This can readily be explained by mass creation very close to the center.

d) Observations using the HST have shown what appear to be very young globular clusters in the central region of the well-known radio galaxy

NGC 1275 (Holtzman et al. 1992). This galaxy gives every indication that parts of it at least have an age $\sim 10^{10}$ yr. Both of these examples suggest that mass creation may well be continuing at a low level in the nuclear regions of old galaxies.

A more general question is to what extent there is good evidence that the majority of galaxies have ages $\sim 10^{10}$ yr. In fact we can only determine ages accurately when we can observe the color-magnitude diagrams of clusters of coeval stars and detect their turn-off on the main sequence or the positions of their horizontal branches. This restricts us so far to our Galaxy and objects like the LMC. The arguments based on color measurements which have been made in general (cf. Larson & Tinsley 1978) involve assumptions about the mix of stellar populations that are not very secure, though they are often implicitly assumed to support the claim that the majority of galaxies are old. If we look at dynamical motions, clusters of galaxies like the Coma Cluster appear from their forms to be totally relaxed, and this means that they are very well mixed. Since in a typical case it takes about 10^9 yr for a galaxy to traverse a cluster diameter, such cluster galaxies should have ages $\gtrsim 10^{10}$ yr. On the other hand, there are many clusters which are clearly not relaxed (e.g., the Hercules Cluster) which contain many bright spirals. The dynamical argument would then suggest that the age is no more than $\sim 10^9$ yr. This was the position taken by Ambartsumian (1958, 1965) when he first discussed expanding associations of galaxies.

So how are galaxies formed in this model? It appears likely that there are two routes. Given a small seed mass, creation in the vicinity of the center will add to that mass which will be driven out but will still remain bound to the overall system. In this way, galaxies will increase in mass as a function of time.

On the other hand, creation at the center will lead to matter which is ejected from the galaxy and can form seeds for the formation of new galaxies.

Explosive ejections of this type is what we see in radio galaxies, QSOs and other active nuclei.

The classical explanation of the latter phenomena is that they arise in a rather mysterious way after some matter from an accretion disk falls into a black hole. A discussion of the many ingenious arguments which have been made has recently been summarized by Blandford & Rees (1992). How are these phenomena to be alternatively explained in a mass-creation event?

The creation units described in this paper have an early stage in which gravitational fields are strong, with creation taking place close to an event horizon, which introduces a time dilatation factor large compared with unity, a factor upward of 10^6 for large creation units. This influences the time scale as measured by an external observer in which the creation unit bursts away from its state near the event horizon, an effect of a strong negative pressure term in the dynamical equations similar to that in the inflationary model (Guth 1981). Now it seems unlikely, especially for a rotating object, that the time dilatation factor will be everywhere the same. Using spherical polar coordinates, the dilatation factor will not be precisely of the form $(1 - 2GM/R)^{-1/2}$ but will also have some dependence on the angular coordinates as in the Kerr metric. To an external observer the moment of breakout from the strong gravitation field will therefore appear dependent on θ and φ. Even though to a comoving observer the times of breakout may not be greatly variable with respect to θ, φ, to an external observer the situation will appear otherwise. That is to say instead of the object expanding after the creation phase as a uniform object, it is likely to emerge in a series of blobs or jerks, every blob appearing as a distinct object in its own right. This type of model may be important in attempts to understand the properties of radio galaxies and other active objects in which matter and high-energy particles are ejected in jets.

Let us consider a few examples starting with M87, the classical radio galaxy with a jet interpreted in this way. The well-known synchrotron jet in the large galaxy M87 lies in position angle 290°. It has been known since 1960 that the position angle of the line joining M87 to the radio galaxy M84 is coincident with the position angle of the jet (Wade 1960). This is as direct evidence as one can have of a changing gravitating object at the center of M87, a blob that later becomes the galaxy which we know as M84. Arguing from the angular coincidence of the position angles, the probability of this being so was already about 100 to 1 in favor, while now we have at least the beginning of a theoretical explanation of how such situations arise. Indeed, taking place repeatedly from the breakup of an initial large mass, the product would be a cluster of galaxies. It is an interesting possibility that the Virgo Cluster was formed in this way or, at any rate, the elliptical galaxies in the Virgo Cluster. The spirals are more likely to have been formed as gas from the object, created as M87, impinging on gas from other neighboring creation units. The circumstances that we still see the jet of M87 suggests that the process of forming M84 occurred fairly recently and that here we have good evidence of a young galaxy.

In addition to this, Arp (1987) has shown that a number of X-ray-emitting QSOs are also aligned in the position angle of the jet, and this also is very suggestive of ejection.

We suggest that it is this type of event – creation process in the center of a massive object – which is largely responsible for the generation of powerful radio sources associated with elliptical galaxies.

There is nearly always a preferred axis of ejection in a powerful source, and many optical and near-IR observations have shown that there is a great deal of optical emission along the major axis of the radio emission. Most of these galaxies are very faint, and many have $z > 1$. In some cases blobs are seen, but most of them are far enough away so that structure of the kind seen in M87 will not be detected.

The recent studies (McCarthy et al. 1991: McCarthy, Elston, & Eisenhardt 1992a; McCarthy, Persson, & West 1992b; McCarthy 1991) show that the optical emission is correlated with the radio emission, and everything suggests that the activity takes place from the inside out.

We would also like to explain the existence of QSOs ejected from galaxies as a variation on this same process, but here much more work needs to be done before we have a satisfactory model. Certainly the ejection process may be understood within the framework of the creation in an active nucleus, and there are some very well aligned ejection cases (cf. Arp 1987; Burbidge et al. 1990). While we believe that the large number of associations between bright galaxies (often spirals) and QSOs with large redshifts provide overwhelming evidence for non-cosmological redshifts, we have not yet found any way of explaining these redshifts using the theory outlined in this paper. We believe that this aspect of the problem of violent activity remains a major challenge.

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