Fermi National Accelerator Laboratory

FERMILAB-Conf-85/80-A
June 1985

Cosmology and Particle Physics: A General Review*

Keith A. Olive

Astrophysics Theory Group
Fermilab
Batavia, IL 60510 U.S.A.

*To be published in the Proceedings of the Twelfth Texas Symposium on Relativistic Astrophysics, ed. by M. Livio, Jerusalem, December 1984.

Operated by Universities Research Association Inc. under contract with the United States Department of Energy

I. INTRODUCTION

At the time of the first Texas Symposium in 1963, the field which I am now about to discuss, did not yet exist. Major topics at that symposium revolved around quasi-stellar sources, massive stars and gravitational collapse. The state of cosmology at that time was also in sharp contrast with our present views. After all, the First Texas Symposium was held before the discovery of the 3°K microwave background radiation by Penzias and Wilson.[1)] Discussions on cosmology still centered on trying to determine a basic cosmological model, the steady-state theory or the big bang. Only since the discovery of the microwave background, have efforts been concentrated on the big bang model, i.e. our Universe has evolved from a once very hot and very dense epoch.

Traced backwards, the 3°K photon background gives us information about the Universe when those photons were last in thermal contact with each other; when the age of the Universe was about 10^5 yrs and its temperature about 10^4°K. Our next single most important piece of evidence regarding our hot past comes from the abundances of the light elements. As early as the 1940's, nuclear physics began to play a role in cosmology when it was realized[2)] that if our Universe began with temperatures in excess of 1 MeV or 10^{10}°K, nucleosynthesis would have occurred and in particular the abundance of ^{4}He can be calculated. The fact that big bang nucleosynthesis agrees with the observed abundances of the light elements gives us added proof of our hot beginnings. Big bang nucleosynthesis began when the age of the Universe was about one second.

Cosmology today does not stop at $t = 1s$. Indeed "reasonable" statements begin at the Planck epoch or when $t_u = 10^{-44}$ s. In this review, I hope to highlight our current understanding of the various stages in the evolution of the Universe from $t \sim 10^{-4}$ s to the period of galaxy formation at $t \gtrsim 10^5$ yrs. I will try to follow a chronological order for the discussion. Therefore, I will begin briefly (as I do not believe that too much can actually be said realistically) with the Planck epoch. In section 3, I will discuss the inflationary epoch. In section 4, I will review the present status of big bang baryosynthesis i.e., the origin of the apparent slight excess of baryons over antibaryons. This is perhaps our third most reliable piece of evidence indicating a hot big bang. I will review the present status of big bang nucleosynthesis in section 5 and discuss why I feel it is one of the greatest successes of the standard big bang model. Finally, in section 6, I will review the present role of particles in the Universe. That is their effects on galaxy formation and constraints from present observations that can be placed on particle properties.

Throughout this paper I will be using units such that $\hbar = c = k_B = 1$ and all masses will be given in GeV (unless specifically noted otherwise).

II. THE PLANCK ERA

The Planck era is defined as the epoch when energy scales become comparable to the Planck mass $M_p = 1.22 \times 10^{19}$ GeV. At these energies, one expects gravitational interactions to become comparable to other particle interactions, $G_N = M_p^{-2}$. The age of the Universe at this point is $t_u = M_p^{-1} \simeq 5 \times 10^{-44}$s. Cosmology at or above M_p, must rely on a deep understanding of gravitational interactions. At the present, such theories are not yet available and cosmological models begin to get very fuzzy (or foamy?[3]). There are however several approaches to attack this problem which I will very briefly describe. These include quantum gravity, Kaluza-Klein theories and supergravity.

Quantum gravity[4] is an attempt to describe gravitational interactions at the same level as is possible for the strong, weak and electromagnetic interactions. In the context of the early Universe, quantum gravitational effects on particle production have been discussed[5] for isotropic as well as anisotropic models. Initial conditions such as the primordial wave function of the Universe[6] have also been put forward in this context.

Kaluza-Klein theories[7] which began in the 1920's, have awoken interest again recently.[8] The basic idea is that one associates gauge interactions with extra dimensions. Or rather, one begins with a $4 + d$ dimensional theory and one tries to reduce it to a 4-dimensional space-time with the compactified dimensions acting like gauge interactions. For example, to account for the U(1) gauge group for electromagnetic interactions one must simply go to 5 dimensions. For a SU(5) grand unified theory (GUT) one needs ≥ 11 dimensions. A major problem concerning these theories is that they contain only real-field representations for particles.[9] The standard low energy theory of electroweak interactions however contains chiral fields (i.e., there are left-handed and right-handed helicity representations). Hence Kaluza-Klein theories could not provide a true unified theory which includes the standard low energy world. These theories have been useful however in that they have begun to establish a formalism for treating (and compactifying) extra dimensions which may be present in the more hopeful supergravity or superstring theories. For a review of the cosmological applications of Kaluza-Klein theories see ref. 10.

Supersymmetry[11)] is a symmetry between bosons and fermions. Besides its beauty as a symmetry of nature, it has gained most of its popularity through its resolution of what is known as the gauge hierarchy problem in standard GUTs. Very simply, the gauge hierarchy problem is the problem concerning mass scales which arise in a theory. For example the weak interaction scale is $M_W \sim 10^2$ GeV while the GUT scale is $M_X \sim 10^{15}$ GeV. The problem is why are these scales so different. Furthermore, a technical problem arises when one considers radiative corrections to these scales. Radiative corrections to the weak mass scale will tend to be as large as the GUT scale and must therefore be cancelled with enormous precision, and to many orders in perturbations theory.

Although such a cancellation is possible, it is not at all natural. Supersymmetry resolves[12)] this difficulty in the sense that one can show that these radiative corrections vanish exactly.[13)] Thus the weak scale of 10^2 GeV is said to be stable with respect to radiative corrections.

The most striking effect of making a model supersymmetric, is that one essentially doubles the number of known particles. To all spin 1 particles such as the photon or gluons one adds spin 1/2 partners called photinos and gluinos; to all spin 1/2 leptons and quarks one adds spin 0 partners, sleptons (such as the selectron) and squarks. Spin 0 Higgs bosons are paired up with spin 1/2 Higgsinos. If supersymmetry is made local (supergravity), then the theory incorporates gravity as well and hence the spin 2 graviton is joined to a spin 3/2 gravitino.

If supersymmetry were an exact symmetry of nature, the supersymmetric partners would have identical properties (except for spin) and hence the selectron mass would be degenerate with the electron mass. However charged spin 0 particles with a mass of 511 keV have not been observed. Hence supersymmetry must be broken. In order to preserve the gauge hierarchy the corrections to scalar masses must be kept as small as 10^2 GeV. These corrections turn out to be proportional to the gravitino mass[14)] which is related to the supersymmetry breaking scale by

$$m_{3/2} = M_S^2/M_p \tag{2.1}$$

where $M_p = 1.2 \times 10^{19}$ GeV is the Planck mass. Thus $M_S \lesssim 10^{10}$ GeV.

Low energy local N=1 supergravity has been studied extensively.[15)] At the Planck scale, there are many questions still to be answered. For example, where did this effective low energy N=1 theory come from. Several possibilities have been put forward. The true theory might be an N=8 supergravity[16)] (the largest possible in four dimensions) which breaks down to an N=1. N>1 supergravity theories however, do not contain chiral fields and must be written in terms of constituent particles (preons) which confine at the Planck scale to ordinary particles.[17)] More recently, N=1 supergravity theories expressed in terms of 10-dimensional strings (superstrings)[18)] have drawn much interest in that they seem to be free of problems concerning infinities which are present in all other theories.

If the reader has not yet determined this, my personal bias lies with supergravity. In the next section, I discuss how supersymmetry is helpful to inflation and in section 6 how supersymmetry may provide an answer to the dark matter problems which plague cosmology.

III. INFLATION

As there are already several reviews[19] about inflation,[20] I will try to be brief here. However any review of the very early Universe would be incomplete if it did not at least touch upon inflation. In short what is meant by inflation, is the effect of exponential expansion due to a supercooled phase transition in order to resolve several finetunings regarding the initial conditions in the standard big bang model.

As examples of these problems, I will briefly describe what is known as the horizon problem and the curvature problem. The horizon volume or causally connected volume today, is just related to the age of the Universe $V_0 \propto t_0^3$. The microwave background radiation with temperature $T_0 \sim 3°K$ has been decoupled from itself since the epoch of recombination at $T_d \sim 10^4 °K$. The horizon volume at that time was $V_d \propto t_d^3$. Now the present horizon volume scaled back to the period of decoupling will be $V_0' = V_0(T_0/T_d)^3$ and the ratio of this volume to the horizon volume at decoupling is

$$\begin{aligned} V_0'/V_d &\sim (V_0/V_d)(T_0/T_d)^3 \\ &\sim (t_0/t_d)^3 \, (T_0/T_d)^3 \sim 10^5, \end{aligned} \tag{3.1}$$

where I have used $t_d \sim 3{\times}10^{12}$ sec and $t_0 \sim 5{\times}10^{17}$ sec. The ratio (3.1) corresponds to the number of regions that were causally disconnected at recombination which grew into our present visible Universe.

The microwave background radiation appears to be highly isotropic. In fact, the limits on the anisotropy put[21])

$$\Delta T/T \lesssim (2\text{-}5) \times 10^{-5} \tag{3.2}$$

This means that on large scales, the Universe must be very isotropic and homogeneous, (any inhomogeneities would also produce fluctuations in the microwave background). The horizon problem, therefore, is the lack of an explanation as to why 10^5 causally disconnected regions at t_d all had the same temperature to within one part in 10^4!

The curvature problem (also known as the flatness or oldness problem) stems from the fact that although the Universe is very old, we still do not know whether it is open or closed. If we look at the Freidmann equation for the expansion of the Universe

$$H^2 = \left(\frac{\dot{R}}{R}\right)^2 = \frac{8\pi G\rho}{3} - \frac{k}{R^2} + \frac{\Lambda}{3} \tag{3.3}$$

where H is the Hubble parameter, R is the Robertson-Walker scale factor, ρ is the total mass energy density, k is the curvature constant ($k=0,\pm1$ for a flat, closed or open Universe) and Λ is the cosmological constant. Neglecting Λ, the curvature term can be expressed in terms of the density parameter

$$\Omega = \rho/\rho_c \tag{3.4}$$

$$\rho_c = 3H_o^2/8\pi G_N \tag{3.5}$$

and the present value of the Hubble parameter, H_o as

$$k/R^2 = (\Omega - 1)\ H_0^2\ . \tag{3.6}$$

If we now use the limits $\Omega < 4$ and $H_0 < 100$ km s^{-1} Mpc^{-1} we can form a dimensionless constant

$$\hat{k} = k/R^2T^2 = (\Omega - 1)\ H_0^2/T^2 \lesssim 3H_0^2/T_0^2 < 2 \times 10^{-58} \tag{3.7}$$

where I have used $T_0 > 2.7°K$. In an adiabatically expanding Universe, $\hat{k}$ is absolutely constant ($R \sim T^{-1}$) and thus the limit (6.4) represents an initial condition which must be imposed so that the Universe will have lived this long looking still so flat.

A more natural initial condition might have been $\hat{k} \sim O(1)$. In this case the Universe would have become curvature dominated at $T \sim 10^{-1}\ M_p$. For $k = +1$, this would signify the onset of recollapse. Even for k as small as $O(10^{-40})$ the Universe would have become curvature dominated when $T \sim 10$ MeV or when the age of the Universe was only $O(10^{-2})$ sec. Thus not only is (3.7) a very tight constraint, it must also be strictly obeyed. Of course, it is also possible that $k = 0$ and the Universe is actually spatially flat.

These are the two main problems that led Guth[20)] to consider inflation. In the problems that were just discussed it was assumed that the Universe has always been expanding adiabatically. During a phase transition, however, this is not necessarily the case. If we look at a potential describing a phase transition from a symmetric false vacuum state $\Sigma = 0$ to the broken true vacuum at $\langle\Sigma\rangle = v$ as in Fig. 1, and we suppose that because of the barrier separating the two minima the phase transition was a supercooled first-order transition. If in addition, the transition takes place at T_c such that $T_c^4 < V_0$, the energy stored in the form of vacuum energy will be released. If released fast enough, it will produce radiation at a temperature $T_R^4 \sim V_0$. In this reheating process entropy has been created and

$$(RT)_f \sim (T_R/T_c)\ (RT)_i \tag{3.8}$$

provided that T_c is not too low. Thus we see that during a phase transition the relation RT ~ constant need not hold true and thus our dimensionless constant $\hat{k}$ may actually not have been constant.

The inflationary Universe scenario,[20)] is based on just such a situation. If during some phase transition, the value of RT changed by a factor of $0(10^{29})$, these two cosmological problems would be solved. The isotropy would in a sense be generated by the immense expansion; one small causal region could get blown up and hence our entire visible Universe would have been at one time in thermal contact. In addition, the parameter $\hat{k}$ could have started out 0(1) and have been driven small by the expansion.

If, in an extreme case, a barrier as in Fig. 1 caused a lot of supercooling such that $T_c^4 << V_0$, the dynamics of the expansion would have greatly changed. In the example of Fig. 1 the energy density of the symmetric vacuum, V_0, acts as a cosmological constant with

$$\Lambda = 8\pi\ V_0/M_p^2. \tag{3.9}$$

If the Universe is trapped inside the false vacuum with $\Sigma = 0$, eventually the energy density due, to say, radiation will fall below the vacuum energy density, $\rho << V_0$. When this happens, the expansion rate will be dominated by the constant V_0 and we will get the De Sitter-type expansion

$$R \sim \exp[Ht], \tag{3.10}$$

where

$$H^2 = \Lambda/3 = 8\pi\ V_0/3M_p^2. \tag{3.11}$$

The cosmological problems could be solved if

$$H\tau \gtrsim 65, \tag{3.12}$$

where τ is the duration of the phase transition and the vacuum energy density was converted to radiation so that the reheated temperature is found by

$$\frac{\pi^2}{30}\ N(T_R)\ T_R^4 = V_0. \tag{3.13}$$

where $N(T_R)$ is the number of degrees of freedom at T_R.

If such a barrier persists down to low temperatures, the phase transition must proceed via the formation of bubbles of the broken phase. The bubble formation rate per unit volume is given by[22)]

$$p \sim Ae^{-B}, \tag{3.14}$$

where $A^{1/4}$ is generally taken to be the overall mass scale in the problem ($A \sim T^4$ or $A \sim M^4$) and B is tunneling action. The transition will take place in such a way so as to minimize the action. The phase transition will be completed when $p > H^4$.

The scenario just described is the original idea of Guth[20] for cosmological inflation. In this scenario, the Universe would undergo a phase transition, say $SU(5) \rightarrow SU(3)\times SU(2)\times U(1)$ in which the potential resembled that in Fig. 1. The Universe would then get hung up in the SU(5) phase down to a very low temperature. After completion of the phase transition, the Universe would reheat to

$$T_R \sim M_X/[N(T_R)]^{1/4}. \tag{3.15}$$

Baryon generation would then follow so long as T_R was not too low. (See next section.)

It is now known that there is a problem with Guth's original idea for inflation.[23] It turns out that the requirement that the Universe supercool for a long time ($H\tau > 65$) is not compatible with $p > H^4$, i.e., the phase transition does not finish. In order to have a long inflationary time scale, a large barrier was necessary so as to be sure that the action for tunneling was also large. It is necessary in this scheme that the initial probability for tunneling be very small. The problem is that under these conditions the tunneling probability never catches up with the expansion rate. As a whole, the Universe remains in the De Sitter state trapped in the symmetric SU(5) vacuum with only a few isolated bubbles containing the true $SU(3)\times SU(2)\times U(1)$ vacuum. Not only is the resulting Universe very inhomogeneous, but each bubble remains empty as all of the energy is stored in the bubble walls and is only released through collisions which in this case do not occur.

The solution to this problem is called the new inflationary Universe[24] and its basic and simple idea is this: tunnel first and inflate later. To realize this type of inflation, one must have a long flat scalar potential. If one can argue (e.g., by thermal effects) that at early times or high temperatures the Universe was in the symmetric phase $\Sigma = 0$ and then at some lower temperature $T \ll T_c$ a bubble is formed. The supercooling may be due to either a barrier as in the previous case or a suppression of thermal fluctuations so that the field Σ rests near the origin. In the case of a barrier, once a bubble is formed, if the potential is very long and flat at values of Σ past the barrier, the potential energy density (approximately constant) will again act like a cosmological constant. If a single bubble were to expand by 29 orders of magnitude, the phase transition need not be completed as in the previous case. The entire visible Universe would be contained within one bubble. The bubble would be filled in this case not by bubble collisions, but by dissipation of the kinetic energy of the scalar field as it finally reaches its global minimum. A generic example of such a potential is shown in Fig. 2.

Popular examples of flat potentials considered for inflation have been the Coleman-Weinberg[25] potentials which are derived by taking first-order radiative corrections to the tree potential. If scalar self couplings are small enough, the tree potential can be neglected and we can concentrate on the corrections. Using an $SU(5)$ Coleman-Weinberg potential has been shown to present several insurmountable difficulties[26,27,28,29]. These range from unnatural finetunings of the self-couplings[27] to transitions occuring to the wrong vacuum state.[29]

The most serious blow to Coleman-Weinberg type inflation comes from the density perturbations which are produced during the rollover.[28] The isotropy of the microwave background radiation tells us that any perturbations produced on large scales must have $\delta\rho/\rho \lesssim O(10^{-4})$. Ideally, what one would want from inflation is what is known as the Harrison-Zeldovich[30] spectrum of density fluctuations. They are also known as scale independent perturbations which are the type most desired for the purposes of galaxy formation. Their magnitude, however, must be $O(10^{-4})$. Any perturbations stronger than this would produce visible anisotropies in the microwave background radiation while weaker perturbations would not have had enough time to grow during the present period of matter domination (since decoupling).

As it turns out, phase transitions, such as the SU(5) transition described above, produce[31] very nearly the Harrison-Zeldovich spectrum which is desired. The perturbations are formed because the field ϕ does not roll down to its global minimum homogeneously. There will, in general, be a time spread over which certain regions roll down faster or slower than others. The density perturbations have been calculated[28] for SU(5) and turn out to be $\delta\rho/\rho \sim 50$, i.e., nearly 5 orders of magnitude too large.

Supersymmetry has led to the resolution of some of these problems.[27,32] In ways similar to those in which supersymmetry resolves the gauge hierarchy problem discussed in the previous section, it relieves the problems of finetuning mass scales in inflation and hence allows for flatter potentials. Another way to make flat potentials is to increase the value of the vacuum expectation value v from $\sim 10^{15}$GeV in GUT models to $v \sim M_p \sim 10^{19}$GeV. This class of models is called primordial inflation.[32] Considering primordial inflation and supersymmetry naturally leads one to inflation in supergravity.[33]

The actual phase transition responsible for inflation is no longer the GUT phase transition and may happen before, during, or after the breaking of GUTs. The scale of supersymmetry breaking M_s must be much lower than the GUT or inflation scale ($M_s \sim 10^{10}$ GeV) in order to preserve the gauge hierarchy. Therefore inflation should not be associated with the supersymmetry breaking transition. Indeed there have been efforts to associate the two, but these models have required large amounts of fine tunings.[34]

The simplest SUSY preserving inflationary models however, run into certain difficulties regarding initial conditions. In other words, if the initial conditions for inflation are determined by high temperature effects[35] it has been shown that these are inconsistent with the requirements for a long inflationary period.[36] Non-minimal supergravity models offer a resolution to this problem.[37] In fact a very simple inflationary model can be written[38] in the context of so-called no-scale non-minimal supergravity.[39] These models are attractive in the sense that only one fundamental scale is put in by hand, namely the Planck scale, and the others are determined through radiative corrections. Such non-minimal models are also thought to stem from extended supergravity models of the type discussed in the previous section, and perhaps recent superstring theories.[40]

Finally, further recent work regarding the role of initial conditions and the possibility of inflation seems to favor heavily models of primordial inflation[41]. It has been claimed[42] that at very high temperatures, the field responsible for inflation will not be found near $\Sigma = 0$ but rather spread very far from the high temperature minimum. This however is true only for certain circumstances and it has been shown that this effect does not occur in models of primordial inflation.

The actual origin of inflation and the exact identity of the phase transition producing the desired supercooling are not known. Inflation exists as a possibility within the context of supergravity models. Its apparent beauty still reaches out further than our apparent ignorance.

IV. GUTs AND COSMOLOGY

The origins of the modern connection between particle physics and cosmology really began with the generation[43] of a small but finite baryon to entropy ratio using grand unified theories (GUTs).[44] The problem in cosmology is basically that there is apparently very little antimatter in the Universe and the number of photons greatly exceeds the number of baryons. If we define

$$\eta = (n_B - n_{\bar{B}})/n_\gamma \qquad (4.1)$$

where $n_{B,\bar{B},\gamma}$ is the number density of baryons, antibaryons and photons, we find that

$$\eta = n_B/n_\gamma \sim 10^{-10} - 10^{-9} \qquad (4.2)$$

(see section V). In a standard model, the entropy density today is related to n_γ by

$$s \approx 7n_\gamma \tag{4.3}$$

so that eq. (4.2) implies $n_B/s \sim 10^{-11} - 10^{-10}$. This ratio is conserved however and hence represents another undesirable initial condition, with its origin unknown.

Let us for the moment, assume that in fact $\eta = 0$. We can compute the final number density of nucleons left over after annihilations have frozen out. At very high temperatures (neglecting a quark-hadron transition) $T > 1$ GeV, nucleons were in thermal equilibrium with the photon background and $n_N = n_{\bar{N}} = 3/2n_\gamma$ (a factor of 2 accounts for neutrons and protons and the factor 3/4 for the difference between fermi and bose statistics). As the temperature fell below m_N, annihilations kept the nucleon density at its equilibrium value $(n_N/n_\gamma) = (m_N/T)^{3/2}$ $\exp(-m_N/T)$ until the annihilation rate $\Gamma_A \approx n_N m_\pi^{-2}$ fell below the expansion rate. This occurred at $T \approx 20$ MeV. However, at this time the nucleon number density has already dropped to

$$n_N/n_\gamma = n_{\bar{N}}/n_\gamma \approx 10^{-18}, \tag{4.4}$$

which is eight orders of magnitude too small[45] aside from the problem of having to separate the baryons from the antibaryons. If any separation did occur at higher temperatures (so that annihilations were as yet incomplete) the maximum distance scale on which separation could occur is the causal scale related to the age of the Universe at that time. At $T \approx 20$ MeV, the age of the Universe was only $t \approx 2\times10^{-3}$ sec. At that time, a causal region (with distance scale defined by 2ct) could only have contained $10^{-5}M_\odot$ which is very far from the galactic mass scales which we are asking for separations to occur, $10^{12}M_\odot$.

A final possibility might be statistical fluctuations, but in a region containing $10^{12}M_\odot$, there are $\sim 10^{80}$ photons so that one would only expect statistical fluctuations to produce an asymmetry $\eta \sim 10^{-40}$! Thus we are left with the problem as to the origin of a small non-zero value for η. We can assume that it was an initial condition to start off with and in a baryon number conserving theory it would remain nearly constant. [The production of entropy (photons) could cause it to fall.] In this case, however, we must still ask ourselves, why is it so small? A more attractive possibility, however, is to suppose that the baryon asymmetry was in some way generated by the microphysics. Indeed, if one can show that a small non-zero value for η developed from $\eta = 0$ (or any other value) as an initial condition, we could consider the question solved. In the rest of this section, we will look at this second possibility for generating a non-zero value of η using GUTs.[44)]

There are three basic ingredients necessary[43)] to generate a non-zero η. They are

1. baryon number violating interactions

2. C and CP violation

3. a departure from thermal equilibrium.

The first condition is rather obvious, unless there is some mechanism for violating baryon number conservation, baryon number will be conserved and an initial condition such as $\eta = 0$ will remain fixed. C and CP violation indicate a direction for the asymmetry. That is, should the baryon number violating interactions produce more baryons than antibaryons? If C or CP were conserved, no such direction would exist and the net baryon number would remain at zero. The final ingredient is necessary in order to insure that not all processes are actually occurring at the same rate. For example, in equilibrium if every process which produced a positive baryon number was accompanied by an equivalent process which destroyed it, again no net baryon number would be produced.

The first two of these ingredients are contained in GUTs, the third in an expanding universe where it is not uncommon that interactions come in and out of equilibrium. In SU(5), the fact that quarks and leptons are in the same multiplets allows for baryon non-conserving interactions such as $e^- + d \leftrightarrow \bar{u} + \bar{u}$, etc., or decays of the supermassive gauge bosons X and Y such as $X \rightarrow e^- + d$, $\bar{u} + \bar{u}$. Although today these interactions are very ineffective because of the masses of the X and Y bosons, in the early Universe when $T > M_X \sim 10^{15}$ GeV these types of interactions should have been very important. C and CP violation is very model dependent. In the minimal SU(5) model, the magnitude of C and CP violation is too small to yield a useful value of η. The C and CP violation in general comes from the interference between tree level and first loop corrections.

The departure from equilibrium is very common in the early Universe when interaction rates cannot keep up with the expansion rate. In fact, the simplest (and most useful) scenario for baryon production makes use of the fact that a single decay rate goes out of equilibrium. It is commonly referred to as the out of equilibrium decay scenario.[46] The basic idea is that the gauge bosons X and Y (or Higgs bosons) may have a lifetime long enough to insure that the inverse decays have already ceased so that the baryon number is produced by their free decays.

More specifically, let us call X, either the gauge boson or Higgs boson, which produces the baryon asymmetry through decays. Let α be its coupling to fermions. For X a gauge boson, α will be the GUT fine structure constant, while for X a Higgs boson, $(4\pi\alpha)^{1/2}$ will be the Yukawa coupling to fermions. The decay rate for X will be

$$\Gamma_D \sim \alpha M_X. \tag{4.5}$$

However decays can only begin occurring when the age of the Universe is longer than the X lifetime Γ_D^{-1}, i.e., when $\Gamma_D > H$

$$\alpha M_X \geq N(T)^{1/2}T^2/M_p \tag{4.6}$$

or at a temperature

$$T^2 < \alpha M_X M_p N(T)^{-1/2}. \tag{4.7}$$

Scatterings on the other hand proceed at a rate

$$\Gamma_s \sim \alpha^2 T^3/M_X^2 \tag{4.8}$$

and hence are not effective at lower temperatures. In equilibrium, therefore, decays must have been effective as T fell below M_X in order to track the equilibrium density of X's (and X's). Thus the condition for equilibrium is that at $T = M_X$, $\Gamma_D > H$ or

$$M_X \lesssim \alpha\, M_p (N(M_X))^{-1/2} \sim 10^{18}\alpha \text{ GeV.} \tag{4.9}$$

In this case, we would expect no net baryon asymmetry to be produced.

For masses $M_X \gtrsim 10^{18}\alpha$ GeV, the lifetime of the X bosons is longer than the age of the Universe when $T \sim M_X$. Decays finally begin to occur when $T < M_X$, however, the density of X's is still comparable to photons $n_X/n_\gamma \sim 1$ whereas the equilibrium density at $T < M_X$ is $n_X/n_\gamma \sim (M_X/T)^{3/2} \times \exp[-M_X/T] << 1$. Hence, the decays are occurring out of equilibrium (inverse decays are not occurring), and we have the possibility for producing a net asymmetry.

Let us now look at what happens during the decay of an $X, \bar{X}$ pair. If we consider the example of the X gauge boson and its decays to $\bar{u}, \bar{u}$ with branching ratio r and net baryon number change $\Delta b_1 = -2/3$ and to e^-, d with branching ratio 1-r, and net baryon number change $\Delta b_2 = +1/3$

$$X \xrightarrow[r]{} \bar{u} + \bar{u} \qquad \Delta b_1 = -2/3 \tag{4.10a}$$

$$X \xrightarrow[1-r]{} e^- + d \qquad X \quad \Delta b_2 = +1/3 \tag{4.10b}$$

A similar set of decays will occur for $\bar{X}$

$$\bar{X} \xrightarrow[\bar{r}]{} u + u \qquad \Delta b_{\bar{1}} = +2/3 \tag{4.11a}$$

$$\bar{X} \xrightarrow[1-\bar{r}]{} e^+ + \bar{d} \qquad \Delta b_{\bar{2}} = -1/3 \tag{4.11b}$$

If C and CP are violated then $r \neq \bar{r}$ and we can define the total net baryon number produced per decay of X and $\bar{X}$

$$\begin{aligned} \Delta B &= (\Delta b_1)r + (\Delta b_2)\ (1-r) + (\Delta b_{\bar{1}})\bar{r} + (\Delta b_{\bar{2}})\ (1-\bar{r}) \\ &= \bar{r} - r. \end{aligned} \tag{4.12}$$

The value of $\bar{r}-r$ will of course depend on the specific model for C and CP violation.

The total baryon density that will have been produced by the X, $\bar{X}$ pair [provided Eq. (4.9) is not satisfied] is

$$n_B = (\Delta B)n_X \tag{4.13}$$

and since we also have $n_X = n_{\bar{X}} = n_\gamma$,

$$n_B = (\Delta B)n_\gamma . \tag{4.14}$$

Although the net baryon number is conserved during the subsequent evolution of the Universe, the photon number density is not. A more useful quantity just after baryon generation is the baryon-to-specific entropy ratio, n_B/s. The entropy density, is

$$s = \frac{2\pi^2}{45}\ N(T)T^3 \tag{4.15}$$

At $T \lesssim M_X \sim 10^{15}$ GeV, we expect $N(T) \gtrsim 0(100)$ so that $s \sim 0(100)\ n_\gamma$. Thus the baryon-to-entropy ratio we would expect to produce in the out-of-equilibrium decay scenario would be

$$n_B/s \sim 10^{-2}(\Delta B). \tag{4.16}$$

The value of n_B/s that we are looking for must be related to the limits on η which will be discussed in section V. η in the range $(3\text{-}10)\times10^{-10}$ corresponds to a value of n_B/s in the range $(4.3\text{-}14)\times10^{-11}$. Comparing this with the expected production, Eq. (4.16) gives us a lot of hope that GUTs may provide us with a viable mechanism for generating a small (but not too small) value for η.

Although we can be encouraged by the above scenario, we must still show that given a GUT, after the full set of Boltzmann equations have been integrated, an acceptable and definite value of η emerges. In particular, most GUTs do satisfy Eq. (4.9), for $\alpha = 1/41$ and $M_X \sim 10^{15}$ GeV decays will be occurring at $T \sim M_X$, but in at best partial equilibrium. Thus the estimate, Eq. (5.14) is not a good one.

In Fig. 3, we look at the typical results which one finds after a complete numerical integration[47] of the Boltzmann equations. These particular results are for an SU(5) model, but their behavior is generic for most any GUT. What is plotted is the time development of the baryon-to-entropy ratio n_B/s normalized to the net baryon number produced by pair decay, ΔB. The horizontal scale, M_X/T, is proportional to $t^{1/2}$. The three curves correspond to different choices for the mass of the boson X. In curve 1, we have chosen, a mass which we expect to satisfy the out-of-equilibrium condition $M_X = 3\times10^{18}\alpha$ and we indeed find that the maximum asymmetry has been generated $n_B/s \approx 10^{-2}\Delta(B)$ as we expected (4.16). This in itself confirms the original idea.

The good news that we find from Fig. 3 is that even for lower masses, an asymmetry is still produced. In curve 2, we have chosen M_X =

$3\times10^{17}\alpha$ and we find still a substantial asymmetry $n_B/s \sim 10^{-4}$ (ΔB). What is happening is that at $T \sim M_X$, inverse decays are still effective in trying to restore equilibrium. Eventually, they too freeze out and any X's and $\bar{X}$'s still present, decay freely to produce a net baryon number. If we continue to lower the mass as in curve 3, $M_X = 3\times10^{16}\alpha$, scatterings begin to play a role in driving things further towards equilibrium. Again, when they freeze out the remaining $X,\bar{X}$ pairs decay leaving an asymmetry. If scatterings become dominant, however, the resulting asymmetry in the standard model will become exponentially small with decreasing M_X as shown in the dashed curve. In Fig. 4 we have plotted the final asymmetry which is produced as a function of $K = 3\times10^{17}\alpha/M_X$ where K is defined by

$$K \equiv \Gamma_D/H|_{T=M_X}. \tag{4.17}$$

Depending on whether or not X is gauge or Higgs boson, the resulting final asymmetry can be approximated by

$$n_B/s \approx 2\times10^{-3}(\Delta B)/[1 + (3K)^{1.2}] \tag{4.18}$$

for Higgs bosons, and

$$n_B/s \approx 8\times10^{-3}(\Delta B)/[1 + (16K)^{1.3}] \tag{4.19}$$

for gauge bosons.

Thus we see that GUTs do indeed offer an explanation to the small but finite baryon to entropy ratio. In supersymmetric theories, the ideas are generally the same although the details may be somewhat different.[48),49)]

V. BIG BANG NUCLEOSYNTHESIS

As was noted in the introduction, the two most important pieces of evidence in support of the standard big bang model are the observation[1)] of the 3°K microwave background radiation and the explanation[2)] of the origin of the light elements and their abundances. Because of the initially high temperatures and densities and the large abundance of neutrons relative to protons, the chains of nuclear reactions similar to those occurring in stars might have occurred. Indeed in the simplest model of nucleosynthesis, one can compute the produced abundances of deuterium, ^{3}He, ^{4}He and ^{7}Li and one finds an amazing degree of agreement with the observed abundances. (The observations which must be compared with the big bang abundances must be from sources where little or no subsequent nucleosynthesis has taken place.) In this section I will review the predictions of big bang nucleosynthesis and its cosmological consequences in terms of limits on particle physics.

The temperature region of interest is one typical of nuclear energies, i.e., $T \sim 1$ MeV. The initial conditions for the problem will therefore be set at $T \gg 1$ MeV. Once again, because the asymmetry between baryons and antibaryons is so small and since we do not expect very different asymmetries among the leptons (standard GUT models even predict their similarity) we will take all chemical potentials to be zero. One of the chief quantities of interest will be the neutron-to-proton ratio (n/p). At very high temperatures ($T \gg 1$ MeV), the weak interaction rates for the processes

$$n + \nu_e \leftrightarrow p + e^-$$
$$n + e^+ \leftrightarrow p + \bar{\nu}_e \qquad (5.1)$$
$$n \leftrightarrow p + e^- + \bar{\nu}_e$$

were all in equilibrium, i.e., $\Gamma_W > H$. Thus we would expect that initially $(n/p) \approx 1$. Actually in equilibrium, the ratio is essentially controlled by the boltzmann factor so that

$$(n/p) \simeq \exp(-\Delta m/T), \qquad (5.2)$$

where $\Delta m = m_n - m_p$ is the neutron-proton mass difference. For $T \gg \Delta m$, $(n/p) \simeq 1$.

At temperatures $T \gg 1$ MeV, nucleosynthesis cannot begin to occur even though the rate for forming the first isotope, deuterium, is sufficiently rapid. To begin with, at $T \gtrsim 1$ MeV deuterium is photodissociated because $E_\gamma > 2.2$ MeV (the binding energy of deuterium; $E_\gamma \approx 2.7T$ for a blackbody). Furthermore, the density of photons is very high $n_\gamma/n_B \sim 10^{10}$. Thus the onset of nucleosynthesis will depend on the quantity

$$\eta^{-1} \exp[-2.2 \text{ MeV}/T] \qquad (5.3)$$

where η is defined as before. When this quantity (5.3) becomes $\lesssim O(1)$, the rate for $p + n \rightarrow D + \gamma$ finally becomes greater than the rate for dissociation $D + \gamma \rightarrow p + n$. This occurs when $T \sim 0.1$ MeV or when the Universe is a little over 2 min. old.

Because nucleosynthesis begins when $T < 1$ MeV, the rates for processes which control (n/p) (5.1) as well as those which keep neutrinos in equilibrium are frozen out. Furthermore, because the rates for processes (5.1) also freeze out (at $T \lesssim 1$ MeV), the neutron to proton ratio must be adjusted from its equilibrium value. When freeze out occurs, the ratio (n/p) is relatively fixed at

$$(n/p) \sim 1/6. \tag{5.4}$$

This equilibrium value is adjusted by taking into account the free neutron decays up until the time at which nucleosynthesis begins. This reduces the ratio to

$$(n/p) \sim 1/7. \tag{5.5}$$

Since virtually all the neutrons available end up in deuterium which gets quickly converted to ^4He, we can estimate the ratio of the ^4He nuclei formed compared with the number of protons left over

$$X_4 \equiv (N_{^4He}/N_H) = 1/2\ (n/p)\bigl(1 - (n/p)\bigr) \tag{5.6}$$

or more importantly the ^4He mass fraction

$$Y_4 = 4X_4/(1 + 4X_4) = 2(n/p)/\bigl(1 + (n/p)\bigr). \tag{5.7}$$

For $(n/p) \simeq 1/7$, we estimate that $Y_4 \simeq 0.25$ which is very close to the observed value.

The actual calculated value of Y_4 will depend on a numerical calculation which runs through the complete sequence of nuclear reactions.[50] The nuclear chain is temporarily halted because there are gaps at masses A = 5 and A = 8, i.e., there are no stable nuclei with those masses. There is some further production, however, which accounts for the abundances of ^{6}Li and ^{7}Li. Once again because of the gap at A = 8 there is very little subsequent nucleosynthesis in the big bang. A second chief factor in the ending of nucleosynthesis is that during this whole process the Universe continues to expand and cool. At lower temperatures it becomes exponentially difficult to overcome the Coulomb barriers in nuclear collisions. In spite of these effects, numerical calculations of the elemental abundance continue the chain up until Al.

Before reviewing the results of the big bang nucleosynthesis[50-53] calculations, it is important to realize that there are three additional parameters which have a very strong effect on the results.* They are 1) the baryon-to-photon ratio η; 2) the neutron half-life $\tau_{1/2}$; 3) the number of light particles or, in particular, the number of neutrino flavors N_ν.

As we have seen above, the value of η controls the onset of nucleosynthesis (5.3). Basically what happens is that for a larger baryon-to-photon ratio η the quantity (5.3) becomes smaller thus allowing nucleosynthesis to begin earlier at a higher temperature. Remember also that a key ingredient in determining the final mass fraction of ^{4}He, Y_4, was (n/p) [see Eq. (5.7)] and that the final value of (n/p) was determined by the time at which nucleosynthesis begins thus controlling the time available for free decays after freeze out. If nucleosynthesis begins earlier, this leaves less time for neutrons to decay and the value of (n/p) and hence Y_4 is increased.

---

*I am not considering the effects of a chemical potential, which can also greatly vary the results.[54]

The value of η cannot be determined directly from observations. If we break e find that

$$n_B = \rho_B/m_B = \Omega_B \rho_c/m_B$$
$$= 1.13 \times 10^{-5} \Omega_B h_0^2 \text{ cm}^{-3}, \tag{5.8}$$

where ρ_B is the energy density in baryons, m_B is the nucleon mass, Ω_B is that part of Ω which is in the form of baryons and ρ_c is the critical energy density. The number density of photons is just

$$n_\gamma = 400\ (T_0/2.7)^3 \text{ cm}^{-3}, \tag{5.9}$$

where T_0 is the present temperature of the microwave background radiation. Putting η back together we find

$$\eta = 2.81 \times 10^{-8} \Omega_B h_0^2 (2.7/T_0)^3. \tag{5.10}$$

Thus we could determine η if we knew Ω_B, h_0, and T_0.

The second parameter, $\tau_{1/2}$, is important in that it also plays a role in determining the value of Y_4. Although we don't usually consider $\tau_{1/2}$ a parameter, the uncertainties in its measured value are significant from the point of view of nucleosynthesis. After all, it is this quantity which will control the weak interaction rates and hence determine the freeze-out temperature. The common value of $\tau_{1/2} = 10.6$ min. is actually uncertain by about two percent and this is enough to affect the production of ^{4}He. The range we will consider is

$$10.4 \text{ min.} < \tau_{1/2} \leq 10.8 \text{ min.} \tag{5.11}$$

As in the case of η, increasing $\tau_{1/2}$ leads to a larger value of Y_4. We can see this by looking again at a comparison between the weak interaction rates and the expansion rate. If we parameterize the weak interaction rate by $\Gamma_{wk} = AT^5$ and the expansion rate by $H = BT^2$ then the freeze-out temperature is determined by

$$H(T_D) = \Gamma_{wk}(T_D) \tag{5.12}$$

or

$$T_d^3 = B/A. \tag{5.13}$$

If we now increase $\tau_{1/2}$, this corresponds to decreasing $\Gamma_{wk} \sim \tau_{1/2}^{-1}$ or decreasing the value of A. This in turn gives a higher value for T_d. Now if T_d is larger, this will give a larger value of (n/p) at freeze-out via Eq. (5.2) and hence more ^{4}He via Eq. (5.7).

The final input parameter, we said was the number of light particles. Specifically, what we mean is the number of degrees of freedom corresponding to particles which are still relativistic ($m \ll T$) when $T < O(1)$ MeV. In addition, we must require that these particles be relatively stable so that they will be present when freeze-out occurs, thus $\tau >$ few seconds. As we hinted to above, likely candidates for these particles are neutrinos and thus the number of neutrino flavors N_ν becomes important. Of course any other types of light particles such as photinos or axions, etc., may also be important.

The number of neutrino flavors N_ν will also affect the primordial abundance of ^4He and like η and $\tau_{1/2}$, increasing N_ν increases Y_4. The expansion rate is proportional to $[N(T)]^{1/2}$. At $T \gtrsim 1$ MeV, N(T) is given by

$$N(T) = 2 + \frac{7}{2} + \frac{7}{4} N_\nu \tag{5.14}$$

which takes into account the contribution of γ's, $e^{\pm}$'s, and N_ν flavors of neutrinos. Thus increasing N_ν, increases B in the notation of Eq. (5.13) and again leads to higher value of T_d, with the same effect of producing more ^4He.

Let us now look at the observations[55] which tells us the abundances of the light elements. In particular, we will be interested in the abundances of D, ^3He, ^4He, and ^7Li. Deuterium is the most easily destroyed of the light elements. It is also very difficult to produce in astrophysical systems where it is not further processed to form ^3He. Therefore, any of the observed D is generally assumed to be primordial. Furthermore because deuterium is so easily destroyed (or burned) we must assume that the abundance of D produced in the big bang is greater than the observed value or

$$(D/H)_{BB} \geq (D/H)_{OBS}, \tag{5.15}$$

where (D/H) is the ratio (by number) of deuterium to hydrogen.

Unlike deuterium, ^3He is very difficult to destroy in its entirety in stellar systems. Pre-main-sequence stars are very efficient in burning deuterium to ^3He via $D + p \rightarrow {}^3He + \gamma$. ^3He is only destroyed at high temperatures ($T > 7 \times 10^6\,°K$) through $^3He + {}^3He \rightarrow {}^4He + 2p$ and $^3He + {}^4He \rightarrow {}^7Be + \gamma$. At higher temperatures ($T > 10^8\,°K$), ^4He is burned to carbon and oxygen. The point is that, in general, some fraction g of the initial ^3He abundance will survive stellar processing. If one takes into account the fact some of this ^3He is redeposited in the interstellar medium (pre-solar) then in terms of g we have

$$(D + {}^3He/H)|_{BB} \leq (D + {}^3He)/H|_{pre\odot} + \left(1/g - 1\right) {}^3He/H|_{pre\odot} \qquad (5.16)$$

The value of g, however, can only be determined[56] by models of stellar evolution and in fact may differ depending on the mass of the star. In low mass stars ($M < 8M_\odot$), $g > 0.7$ is not unreasonable while for high mass stars ($8M_\odot < M < 100\,M_\odot$), g may be as low as 1/4. Since an initial spectrum of stellar masses would cover all ranges, perhaps a lower limit to g of 1/2 - 1/4 would be safe.

Using the observational limits on D/H and $^3He/H$ (see refs. 55,57-60)

$$(D/H) \geq (1\text{-}2) \times 10^{-5} \qquad (5.17a)$$

$$(D+{}^3He)/H|_{pre\odot} \leq 4 \times 10^{-5} \qquad (5.17b)$$

$$^3He/H|_{pre\odot} \leq 2 \times 10^{-5} \qquad (5.17c)$$

we find from the results of big bang nucleosynthesis calculations[53] shown in Fig. 5 that

$$(3\text{-}4) \times 10^{-10} \leq \eta \leq (7\text{-}10) \times 10^{-10} \tag{5.18}$$

to be consistent with both ^{3}He and D.

^{7}Li is another isotope which is in principle difficult to draw solid conclusions from. The main difficulty is that ^{7}Li is both easily produced as well as destroyed. Recently, however, there have been some measurements[61)] of the ^{7}Li abundance in some very old Population II stars. Since some ^{7}Li might have been destroyed before the formation of these stars, we might expect $(^7\mathrm{Li/H})_{\mathrm{PopII}} \leq (^7\mathrm{Li/H})_{\mathrm{BB}}$. (The present ^{7}Li abundance would be larger still representing the contribution from stellar processing.) The observed limit on the ^{7}Li abundance is

$$(^7\mathrm{Li/H})_{\mathrm{PopII}} \lesssim 1.5 \times 10^{-10} \tag{5.19}$$

and is consistent with big bang nucleosynthesis for

$$10^{-10} < \eta < 7 \times 10^{-10} \tag{5.20}$$

which agrees well with (5.18).

This brings us to ^{4}He which is probably the most important of the isotopes studied. The main reason ^{4}He is so important is that there is so much of it. Next to hydrogen it is the most abundant element around and its abundance is quite well known. Unlike the other light elements which have observational uncertainties of $\geq$ 100%, the ^{4}He abundances are measured to within a few per cent. The main problem is that it is also produced in stars and care must be taken in trying to derive the "observed" primordial abundance.

To be sure, one can place an upper limit on the primordial abundance by $Y_{4BB} < Y_{4OBS}$ (Y_4, remember is the total ^{4}He mass fraction). However, in order to use big bang nucleosynthesis to set limits on particle physics (e.g., N_ν) a much more accurate determination of Y_{4BB} is needed. Spectral measurements[55] of galactic HII regions give very accurate values of Y_4, however, there they have been contaminated with by-products of stellar processing. The observations of galaxies with low metal abundances could in principle yield an accurate value of Y_{4BB} but these measurements are difficult because these galaxies are typically very far away. It is not possible within the scope of these lectures to cover completely the discussion of Y_4. The best estimates consistent with the observations place Y_4 in the range

$$0.22 \leq Y_4 \leq 0.25. \tag{5.21}$$

If we restrict ourselves as before to $N_\nu = 3$, $\tau_{1/2} = 10.6$ min., the upper limit on Y_4 implies an upper limit on η from Fig. 6

$$\eta \leq 5 \times 10^{-5} \tag{5.22}$$

which is once again consistent with the previous limits Eq. (5.18). (The lower limit on Y_4 does not give an interesting bound on η.)

Figure 6 actually contains significantly more information than just a limit on η. In Fig. 6, we see clearly the behavior of Y_4 with respect to all three parameters: η, $\tau_{1/2}$, and N_ν. It is clear how Y_4 increases with increasing values of any of the three parameters. It is also immediately clear that we can set a limit [51-53] on N_ν provided that we have a lower limit to η. Using $\eta > 3 \times 10^{-10}$ and $Y_4 < 0.25$, we find that $N_\nu \leq 4$ with the equality being at best marginal. This implies that at most one more generation is allowed, assuming that the neutrinos associated with each generation are light and stable.

The strong dependence of Y_4 on the three parameters requires great precision to strengthen the limits due to nucleosynthesis. Strictly speaking, $\eta > 3 \times 10^{-10}$ and $\tau_{1/2} > 10.4$ min. allows $N_\nu = 4$ only if $Y_4 \geq 0.253$; however, we are not yet in a position to believe the third decimal place. For $\tau_{1/2} \gtrsim 10.4$ min., the limit Eq. (5.22) on η can be relaxed so that $N_\nu \leq 3$, $Y_4 < 0.25$ implies $\eta < 7 \times 10^{-10}$. We can also turn the limits around and set a lower limit to the helium abundance by assuming $\eta > 3 \times 10^{-10}$ and $N_\nu \geq 3$ then we have $Y_4 > 0.24$. If future observations actually yield $Y_4 < 0.24$, one would have to argue that perhaps ν_τ is heavy and unstable (the present limit is only $m_{\nu_\tau} < 250$ MeV). If we only assume $N_\nu \geq 2$, then the lower limit on Y_4 becomes $Y_4 \geq 0.22$. Any observation of the primordial helium abundance less than 0.22 would indicate an inconsistency with the standard model.

There is still one more important consequence of the above limits, that is the limit on η can be converted to a limit on the baryon density and Ω_B. If we turn around Eq. (5.10) we have

$$\Omega_B = 3.56 \times 10^7 \eta h_0^{-2} (T_0/2.7)^3, \qquad (5.23)$$

and using the limits on η Eq. (5.18), h_0 and T_0 from (2.7 - 3)°K we find a range for Ω_B

$$0.01 \leq \Omega_B \leq 0.19. \qquad (5.24)$$

Recall that for a closed Universe $\Omega > 1$, thus from Eq. (5.24) we can conclude that the Universe is not closed by baryons. This does not exclude the possibility that other forms of matter (e.g., massive neutrinos, etc.) exist in large quantities to provide for a large Ω. In fact, if large clusters of galaxies were representative of Ω the limit from nucleosynthesis would indicate that some form of dark matter must exist.

VI DARK MATTER AND GALAXY FORMATION

Inflation predicts that $\Omega=1$. Big bang nucleosynthesis requires $\Omega_B \lesssim 0.2$. The obvious resolution to this conflict is to suppose that there exists non-baryonic dark matter $\Omega_D \gtrsim 0.8$. Observational determinations of Ω however always prove to be less than a few tenths[62] indicating that perhaps a large fraction of the dark matter is unclustered.

The above represents the first part of a growing chain of dark matter problems. Other dark matter problems become evident when one considers smaller scales. On the scale of galaxies, the observation of flat rotation curves implies strongly the existence of a dark halo component for spiral galaxies. That is, it appears that a substantial fraction of the galactic mass is non-visible. Because on galactic[63] scales $\Omega \sim 0.1$, one might think that baryons are the logical candidate. However unless the baryons are carefully hidden in clouds distributed over the halo or in low mass stars with a special mass distribution, baryons are unlikely[64] to provide the missing mass in galactic halos. (Very massive black holes, $M \geq 100\ M_{\odot}$, which leave little or no ejecta remain a possibility).

On smaller scales, those of dwarf spheroidal galaxies, there appears to be a dark matter problem as well.[65] Even locally, in the solar neighborhood, there appears to be as much dark material as is in the form of stars and gas.[66] The dark matter in the galactic disk however is presumably in the form of baryons as they must have undergone dissipational processes to get them into the disk.

Once we accept that non-baryonic dark matter is needed, we can distinguish[67] three forms of dark matter depending on their impact on the growth of density perturbations needed to account for galaxy formation. Perturbations grow primarily during the matter dominated phase of expansion. Particles which are relativistic just before matter dominance are called hot particles. Neutrinos or light Higgsinos are examples. Because these particles are relativistic at relatively late times, their free streaming wipes out perturbations out to scales[68]

$$M_J = 3 \times 10^{18} M_\odot / m_\nu^2(\mathrm{ev}) \tag{6.1}$$

where M_J is the Jeans mass and is the minimum scale on which clustering occurs. In the hot scenario of galaxy formation, large scale structures form first and must fragment down to galactic scales.[69] A second class of dark matter candidates is referred to as warm matter and includes more massive particles up to O(1) keV. In this case, initial mass scales are somewhat smaller.[70,71] If they exist, right-handed neutrinos might be warm particle candidates.[70]

Particles which have been non-relativistic long before perturbation growth began are called cold particles. Particles more massive than 1 keV such as heavy neutrinos, photinos, higgsinos, sneutrinos, etc.

are cold particle candidates. In the cold dark matter scenario[72], small scale structures ($M > 10^6 M_\odot$) form first and larger scales are built in a hierarchical manner. For a summary of the pros and cons of these three possibilities in relation to models of galaxy formation, see ref. 73 in these proceedings.

The existence of the dark matter candidate depends of course on the particular particle physics model that one employs. Supersymmetric models are interesting in this context because they guarantee one stable and probably massive particle. If the particle is neutral there are basically four possibilities: 1) the spin 0 partner of the neutrino or sneutrino; 2) the spin 1/2 partner of an axion like particle or axino; 3) the spin 3/2 gravitino and 4) the spin 1/2 partner of the photon or photino.

Similar to limits on neutrino masses, there are limits on most of the supersymmetric dark matter candidates. The exception is the sneutrino, whose mass is not constrained by cosmology[74], but should be ~2 GeV if it accounts for $\Omega=1$. The axino mass limits depend on the temperature at which axino interactions dropped out of equilibrium but roughly $m_{\tilde{a}} \lesssim 1$ keV. Similarly, the gravitino if stable must have[75] $m_{3/2} \lesssim 1$ keV unless the number of gravitinos was reduced by inflation[76] in which case $m_{3/2}$ is unconstrained. Photinos are most like neutrinos in the sense their mass must be[77,78] $m_{\tilde{\gamma}} \gtrsim 1/2$ GeV. This last possibility has been suggested[79] to account for low energy antiprotons observed in cosmic rays through photino annihilations and photinos in the sun may provide[80] a direct observational test of dark matter through annihilations to energetic neutrinos.

The gravitino is also quite interesting cosmologically in that if it is not stable it must be very long lived. Its decay must be purely gravitational and hence its decay rate given by[81]

$$\Gamma \sim m_{3/2}^3/M_P^2 \tag{6.2}$$

so that decays may have been occurring fairly recently. These late decays may be of interest[82] to galaxy formation models using a decaying particle scenario.[83,82] Decays into photons and photinos can highly constrain the abundance of gravitinos[84] again necessitating an inflationary solution[76] but with a low reheat temperature so as not to reproduce gravitinos after inflation. Such low abundances of gravitinos decaying into $\gamma + \tilde{\gamma}$ however may be able to account for observed features in the γ-ray spectrum.[85]

In this last section, we have seen some of the recent results and activity taking place between particle physics and cosmology, and will probably remain the most active area in the interface in the near future. Planck era cosmology must still rely heavily on breakthroughs in particle physics and quantum gravity. Inflation seems to have reached a plateau. Inflation is possible, but its origin is unclear. Perhaps some new idea from superstring theory will replace it. In sections IV and V, big bang nucleosynthesis and baryosynthesis have laid the corner stone of modern cosmology. This leaves us with the dark matter problems, their solutions and tests of the existence of dark matter. The detection[79,80,86] of dark matter offers cosmologists as well as particle physics with a promising challenge.

REFERENCES

1) A. A. Penzias and R. W. Wilson, Ap. J. 142 (1965) 419.

2) G. Gamow, Phys. Rev. 70 (1946) 572; R. A. Alpher, H. Bethe and G. Gamow, Phys. Rev. 73 (1948) 803.

3) S. W. Hawking, Nucl. Phys. B144 (1978) 349.

4) For reviews see articles by B. S. DeWitt and S. W. Hawking in General Relativity: An Einstein Centenary Survey, S. W. Hawking and W. Israel (eds) (Cambridge Univ. Press, 1979).

5) M. V. Fischetti, J. B. Hartle and B.-L. Hu, Phys. Rev. D20 (1979) 1757; J. B. Hartle and B.-L. Hu, Phys. Rev. D20 (1979) 1772; D21 (1980) 2756; J. B. Hartle, Phys. Rev. D22 (1980) 2091; D23 (1981) 2121.

6) J. B. Hartle and S. W. Hawking, Phys. Rev. D28 (1983) 2960.

7) T. Kaluza, Sitzunger Preuss. Akad. Wiss. Math. Phys. K1 (1921) 966; O. Klein, Z. Phys. 37 (1926) 895.

8) E. Witten, Nucl. Phys. B186 (1981) 412; A. Salam and J. Strathdee, Ann. Phys. 141 (1982) 316.

9) E. Witten, in Proc. June 1983 Shelter Island II Conf., eds. N. Khuri etal. (MIT Press, Cambridge MA 1984).

10) E. Kolb, Fermilab preprint 85-17 (1985).

11) Y. A. Gol'Fand and E. P. Likhtman, Pis'ma Zh. Eksp. Teor. Fiz. 13 (1971) 323; D. Volkov and V. P. Akulov, Phys. Lett. 46B (1973) 109; J. Wess and B. Zumino, Nucl. Phys. B70 (1974) 39; for a review see: P. Fayet and S. Ferrara, Phys. Rep. 32C (1977) 249.

12) L. Maiani, in Proc. of the Summer School of Gif-sur-Yvette (1979) p. 3; E. Witten, Nucl. Phys. B188 (1981) 513; S. Dimopoulos and H. Georgi, Nucl. Phys. B193 (1981) 150; N. Sakai, Z. Phys. C11 (1982) 153.

13) J. Wess and B. Zumino, Phys. Lett. 49B (1974) 52; J. Iliopoulos and B. Zumino, Nucl. Phys. B76 (1974) 310; S. Ferrara, J. Iliopoulos and B. Zumino, Nucl. Phys. B77 (1974) 413; M. T. Grisaru, W. Siegel and M. Rocek, Nucl. Phys. B159 (1979) 420.

14) J. Ellis and D. V. Nanopoulos, Phys. Lett. 116B (1982) 133.

15) See reviews by: R. Arnowitt, A. H. Chamseddine and P. Nath, Northeastern Univ. preprints 2597, 2600, 2613 (1983); J. Ellis, CERN preprint TH-3718 (1983); D. V. Nanopoulos, CERN preprint TH-3699 (1983); H.-P. Nilles, Phys. Rep. 110C (1984) 1; J. Polchinski, Harvard Univ. preprint. HUTP-83/A036 (1983).

16) E. Cremmer and B. Julia, Phys. Lett. 80B (1978) 48 and Nucl. Phys. B159 (1979) 141; B. deWit and H. Nicolai, Phys. Lett. 108B (1981) 285 and Nucl. Phys. B208 (1982) 323.

17) J. Ellis, M. K. Gaillard and B. Zumino, Phys. Lett. 94B (1980) 343; J. Ellis, M. K. Gaillard, L. Maiani and B. Zumino, in Unification of the Fundamental Particle Interactions, ed. S. Ferrara, J. Ellis and P. Van Nieuwenhuizen (Plenum Press, NY, 1980) p. 69.

18) J. H. Schwarz, Phys. Rep. 89 (1982) 223; M. B. Green and J. H. Schwarz, Caltech preprints CALT-68-1182, CALT-68-1194 (1984).

19) A. D. Linde, Rep. Prog. Phys. 47 (1984) 925; R. Brandenberger, Rev. Mod. Phys. 57 (1985) 1.

20) A. H. Guth, Phys. Rev. D23 (1981) 347.

21) J. Uson and D. Wilkenson, in Proc. of the Inner Space/Outer Space Conf., eds. E. Kolb, M. S. Turner, D. Lindley, K. A. Olive, and D. Seckel (Univ. of Chicago Press 1985).

22) S. Coleman, Phys. Rev. D15, 2929 (1977); C. Callan and S. Coleman, Phys. Rev. D16, 1762 (1977).

23) A. H. Guth and E. Weinberg, Phys. Rev. D23, 826 (1981); Nucl. Phys. B212, 321 (1983).

24) A. D. Linde, Phys. Lett. 108B, 389 (1982); A. Albrecht and P. J. Steinhardt, Phys. Rev. Lett. 48, 1220 (1982).

25) S. Coleman and S. Weinberg, Phys. Rev. D7, 1888 (1973).

26) A. D. Linde, Phys. Lett. 116B, 335 (1982).

27) J. Ellis, D. V. Nanopoulos, K. A. Olive, and K. Tamvakis, Phys. Lett. 118B, 335 (1982).

28) S. W. Hawking, Phys. Lett. 115B, 295 (1982); A. H. Guth and S.-Y. Pi, Phys. Rev. Lett. 49, 1110 (1982); A. A. Starobinski, Phys. Lett. 117B, 175 (1982); J. M. Bardeen, P. J. Steinhardt, and M. S. Turner, Phys. Rev. D28, 679 (1983).

29) J. Breit, S. Gupta, and A. Zaks, Phys. Rev. Lett. 51. 1007 (1983).

30) E. R. Harrison, Phys. Rev. D1, 2726 (1970); Ya. B. Zel'dovich, Mon. Not. R. Ast. Soc. 160, 1P (1972).

31) W. H. Press, Phys. Scr. 21, 702 (1980).

32) J. Ellis, D. V. Nanopoulos, K. A. Olive, and K. Tamvakis, Nucl. Phys. B221, 224 (1983).

33) D. V. Nanopoulos, K. A. Olive, M. Srednicki, and K. Tamvakis, Phys. Lett. 123B, 41 (1983).

34) B. Ovrut and P. J. Steinhardt, Phys. Rev. Lett. 53 (1984) 732; P. Binetruy and S. Mahajan, LBL preprint 1985.

35) G. Gelmini, D. V. Nanopoulos and K. A. Olive, Phys. Lett. 131B (1983) 53.

36) B. A. Ovrut and P. J. Steinhardt, Phys. Lett. 133B (1983) 161.

37) L. Jensen and K. A. Olive, Fermilab preprint 85-53 (1985).

38) J. Ellis, E. Enqvist, D. V. Nanopoulos, K. A. Olive and M. Srednicki, Phys. Lett. 152B (1985) 175.

39) E. Cremmer, S. Ferrara, C. Kounnas, and D. V. Nanopoulos, Phys. Lett. 133B (1983) 287; J. Ellis, A. B. Lahanas, D. V. Nanopoulos, and K. Tamvakis, Phys. Lett. 134B (1984) 29; J. Ellis, C. Kounnas, and D. V. Nanopoulos, Nucl. Phys. B241 (1984) 406 and Nucl. Phys. B247 (1984) 373.

40) E. Witten, Princeton Univ. preprint February 1985.

41) A. Albrecht and R. Brandenberger, Santa Barbara preprint NSF-ITP-84-146 (1984); G. D. Coughlan and G. G. Ross, Oxford preprint (1984); L. Jensen and K. A. Olive, Fermilab preprint 85/71, 1985.

42) G. Mazenko, W. Unruh, and R. Wald, Phys. Rev. D31 (1985) 273.

43) A. D. Sakharov, Zh. Eksp. Teor. Fiz. Pisma. Red. 5 (1967) 32.

44) For a review see: E. Kolb and M. S. Turner, Ann. Rev. Nucl. Part. Sci. 33 (1983) 645.

45) G. Steigman, Ann. Rev. Astron. and Astrophys. 14 (1976) 339.

46) S. Weinberg, Phys. Rev. Lett. 42, 850 (1979); D. Toussaint, S. B. Treiman, F. Wilczek, and A. Zee, Phys. Rev. D19, 1036 (1979).

47) E. W. Kolb and S. Wolfram, Phys. Lett. B91, 217 (1980); E. W. Kolb and S. Wolfram, Nucl. Phys. B172, 224 (1980); J. N. Fry, K. A. Olive, and M. S. Turner, Phys. Rev. D22, 2953, 2977 (1980).

48) D. V. Nanopoulos and K. Tamvakis, Phys. Lett. 114B (1982) 235.

49) A. Masiero, D. V. Nanopoulos, K. Tamvakis, and T. Yanagida, Z. Phys. C17 (1983) 33.

50) R. V. Wagoner, W. A. Fowler, and F. Hoyle, Ap. J. 148, 3 (1967); R. V. Wagoner, Ap. J. Supp. 18, 247 (1969); R. V. Wagoner, Ap. J. 179, 343 (1973); D. N. Schramm and R. V. Wagoner, Ann. Rev. Nucl. Part. Sci. 27, 37 (1977).

51) G. Steigman, D. N. Schramm, and J. E. Gunn, Phys. Lett. B66, 202 (1977); J. Yang, D. N. Schramm, G. Steigman, and R. T. Rood, Ap. J. 227, 697 (1979).

52) K. A. Olive, D. N. Schramm, G. Steigman, M. S. Turner, and J. Yang, Ap. J. 246, 547 (1981).

53) J. Yang, M. S. Turner, G. Steigman, D. N. Schramm, and K. A. Olive, Ap. J. 281, 493 (1984).

54) Y. David and H. Reeves, in Physical Cosmology, eds. R. Balian, J. Audouze, and D. N. Schramm (North-Holland Pub. Co. Amsterdam, 1980) 443.

55) For a recent compilation of observations see: the Proceedings of the ESO Workshop on Primordial Helium, Eds. P. A. Shaver, D. Kunth, and K. Kajär, Garching, Germany, 1983.

56) I. Iben, Ap. J. 147, 624 (1967); R. T. Rood, Ap. J. 177, 681 (1972); I. Iben and J. W. Truran, Ap. J. 220, 980 (1978); D. S. P. Dearborn, J. B. Blake, K. L. Hainebach, and D. N. Schramm, Ap. J. 223, 552 (1978); W. Brunish and J. W. Truran (in preparation), 1983.

57) D. C. Black, Nature Phys. Sci. 234, 148 (1971); D. C. Black, Geochim. Cosmochim. 36, 347 (1972); J. Geiss and H. Reeves, Astron. Astrophys. 18, 126 (1972).

58) J. T. Trauger etal., Ap. J. Lett. 184, L137 (1973); V. Kunde etal., Ap. J. 263, 443 (1982).

59) D. G. York and J. B. Rogerson, Jr., Ap. J. 203, 378 (1976); A. Vidal-Madjar etal., Ap. J. 211, 91 (1977); A. Vidal-Madjar etal., Astron. Astrophys. 120, 58 (1983).

60) E. Anders, D. Heymann, and E. Mazor, Geochim. Cosmochim. 34, 127 (1970).

61) F. Spite and M. Spite, Astron. Astrophys. 115, 357 (1982); M. Spite and F. Spite, Nature 297, 483 (1982).

62) M. Davis and P. J. E. Peebles, Ann. Rev. Astron. Astrophys. 21 (1983) 109.

63) S. M. Faber and J. J. Gallagher, Ann. Rev. Astron. Astrophys. 17 (1979) 135.

64) D. Hegyi and K. A. Olive, Phys. Lett. 126B (1983) 28 and Fermilab preprint 85/26 (1985).

65) M. Aaronson, Ap. J. 266 (1983) L11; S. M. Faber and D. N. C. Lin, Ap. J. 266 (1983) L17.

66) J. Bahcall, Ap. J. 276 (1984) 169.

67) J. R. Bond and A. S. Szalay, Ap. J. 274 (1983) 443.

68) J. R. Bond, G. Efstathiou and J. Silk, Phys. Rev. Lett. 45 (1980) 1980; Ya. B. Zeldovich and R. A. Sunyaev, Pisma Ast. Zh. 6 (1980) 451.

69) A. Melott, Mon. Not. R. Ast. Soc. 202 (1983) 595; J. R. Bond, A. S. Szalay, and S. D. M. White, Nature 301 (1983) 584.

70) K. A. Olive and M. S. Turner, Phys. Rev. D25 (1982) 213.

71) J. R. Bond, A. S. Szalay, and M. S. Turner, Phys. Rev. Lett. 48 (1982) 1636.

72) P. J. E. Peebles, Ap. J. 263 (1982) L1 and Ap. J. 277 (1984) 440; G. Blumenthal, S. M. Faber, J. R. Primack and M. Rees, Nature 311 (1984) 517.

73) S. D. M. White, these proceedings.

74) L. F. Ibanez, Phys. Lett. 137B (1984) 160; J. Hagelin, G. L. Kane, and S. Raby, Nucl. Phys. B241 (1984) 638.

75) H. R. Pagels and J. R. Primack, Phys. Rev. Lett. 48 (1982) 223.

76) J. Ellis, A. D. Linde, and D. V. Nanopoulos, Phys. Lett. 118B (1982) 39.

77) S. Weinberg, Phys. Rev. Lett. 50 (1983) 387.

78) H. Goldberg, Phys. Rev. Lett. 50 (1983) 1419; L. Krauss, Nucl. Phys. B227 (1983) 556; J. Ellis, J. Hagelin, D. V. Nanopoulos, K. A. Olive, and M. Srednicki, Nucl. Phys. B238 (1984) 453.

79) J. Silk and M. Srednicki, Phys. Rev. Lett. 53 (1984) 624.

80) J. Silk, K. A. Olive, and M. Srednicki, Fermilab preprint 85-62 (1985).

81) S. Weinberg, Phys. Rev. Lett. 48 (1982) 1303.

82) K. A. Olive, D. Seckel, and E. Vishniac, Ap. J. (in press) 1985; K. A. Olive, D. N. Schramm and M. Srednicki, Nucl. Phys. B255 (1985) 495.

83) M. Davis, M. Lecar, C. Pryor, and E. Witten, Ap. J. 250 (1981) 423; P. Hut and S. D. M. White, Nature 310 (1984) 637; M. S. Turner, G. Steigman, and L. Krauss, Phys. Rev. Lett. 52 (1984) 2090; G. Gelmini, D. N. Schramm, and J. W. E. Valle, Phys. Lett. 146B (1984) 311.

84) D. V. Nanopoulos, K. A. Olive and M. Srednicki, Phys. Lett. 127B (1983) 30; M. Yu Khlopov and A. D. Linde, Phys. Lett. 138B (1984) 265; J. Ellis, J. Kim and D. V. Nanopoulos, Phys. Lett. 145B (1984) 181; J. Ellis, D. V. Nanopoulos, and S. Sarkar, CERN preprint Th. 4057 (1985); R. Juszkiewicz, J. Silk, and A. Stebbins, Phys. Lett. B (submitted) 1985.

85) K. A. Olive and J. Silk, Fermilab preprint 85-43 1985.

86) A. Drukier and L. Stodolsky, Phys. Rev. D30 (1984) 2295; M. Goodman and E. Witten, Princeton Univ. preprint (1984).

FIGURE CAPTIONS

Figure 1: The scalar potential for a first order phase transition.

Figure 2: A schematic view of the type of scalar potential needed for new inflation.

Figure 3: The time evolution of the baryon asymmetry in units of (ΔB) for 1) $M_X = 3\times10^{18}\alpha$; 2) $3\times10^{17}\alpha$; 3) $3\times10^{16}\alpha$ and 4) if scatterings remain very effective.

Figure 4: The final baryon asymmetry as a function of $K = 3\times10^{17}\alpha/M_X$ in units of (ΔB). The dashed curve assumes effective scatterings.

Figure 5: The abundances (by number relative to hydrogen) of D, ^3He and their sum as a function of η for $N_\nu = 3$ and $\tau_{1/2} = 10.6$ min.

Figure 6: The abundance (by mass) of ^4He as a function of η for $N_\nu = 2$, 3 and 4, and for $\tau_{1/2} = 10.4$ min. (solid), 10.6 min (dashed), and 10.8 min (dotted).

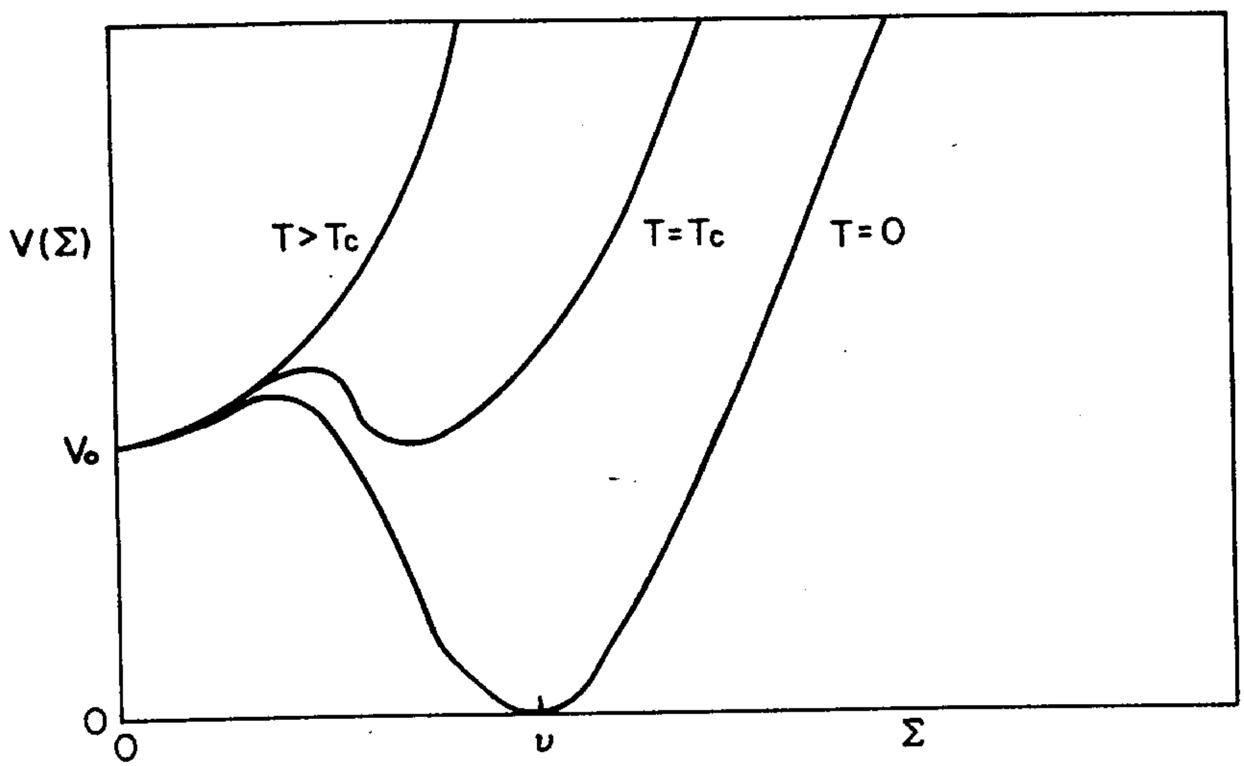

FIGURE 1

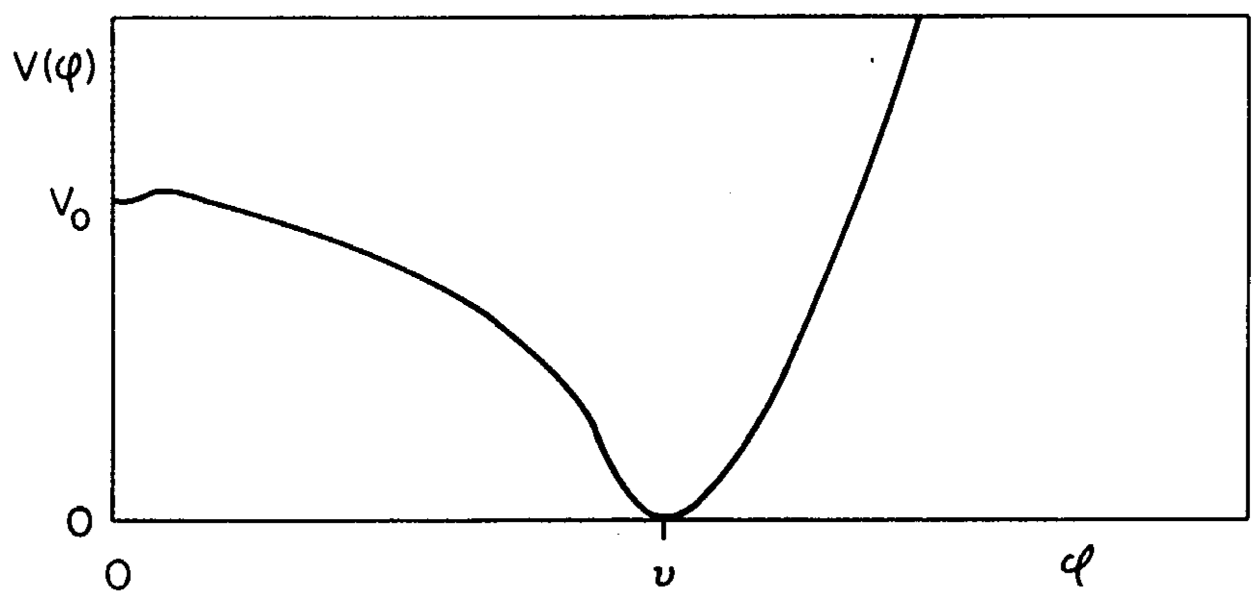

FIGURE 2

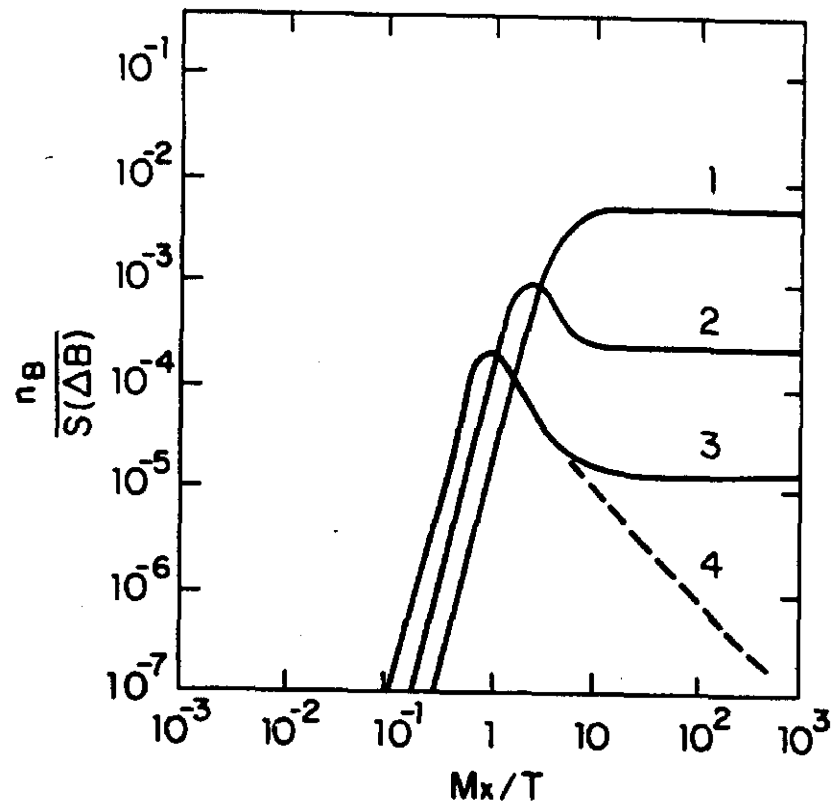

FIGURE 3

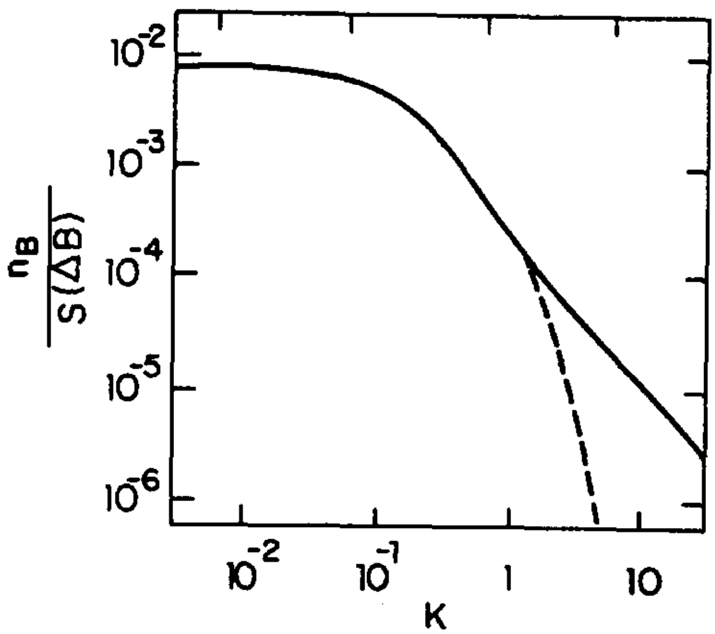

FIGURE 4

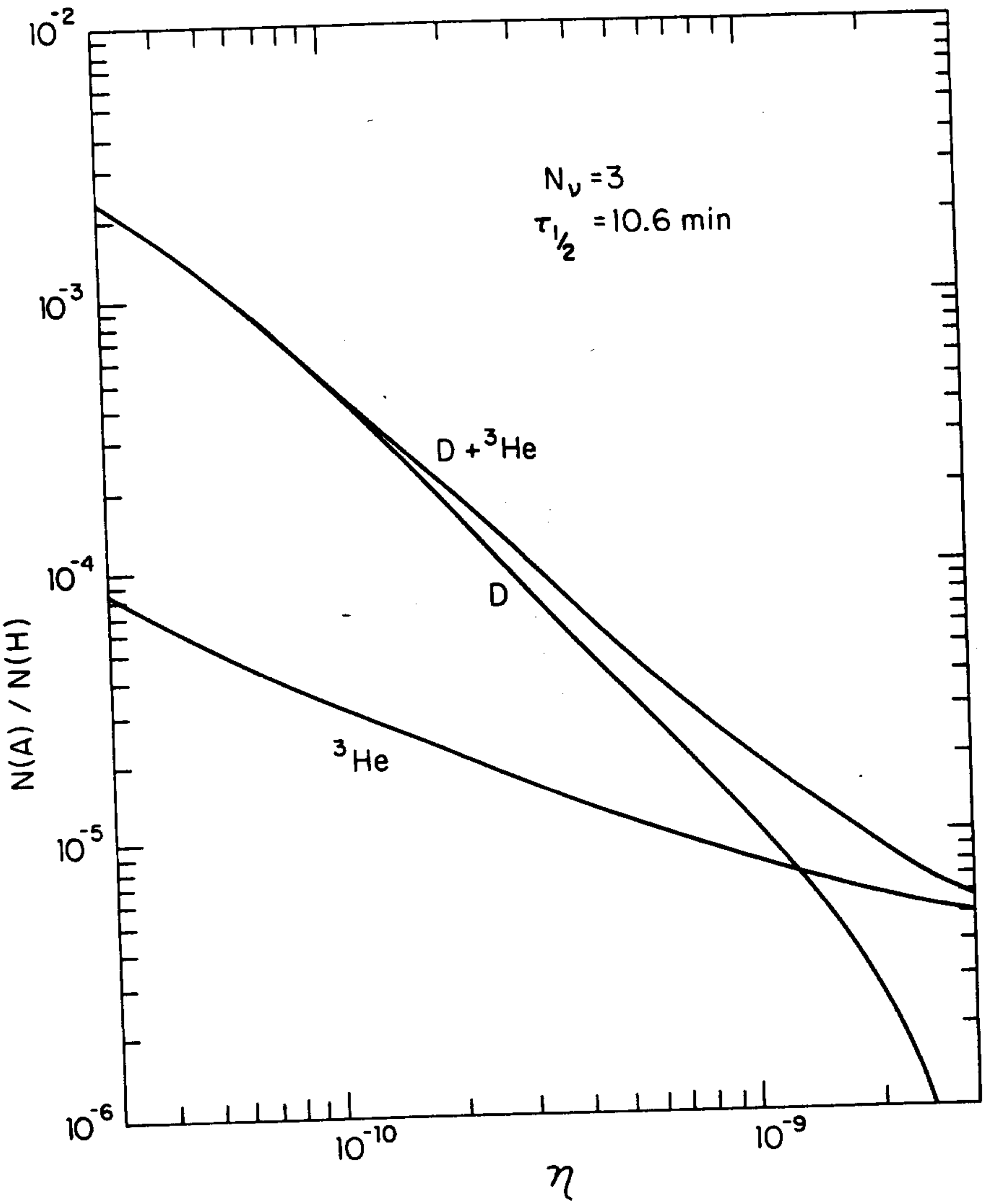
10^{-2}
10^{-3}
10^{-4}
10^{-5}
10^{-6}
10^{-10}
10^{-9}
$N_\nu = 3$
$\tau_{1/2} = 10.6$ min
D + ^{3}He
D
^{3}He
N(A) / N(H)
η

FIGURE 5

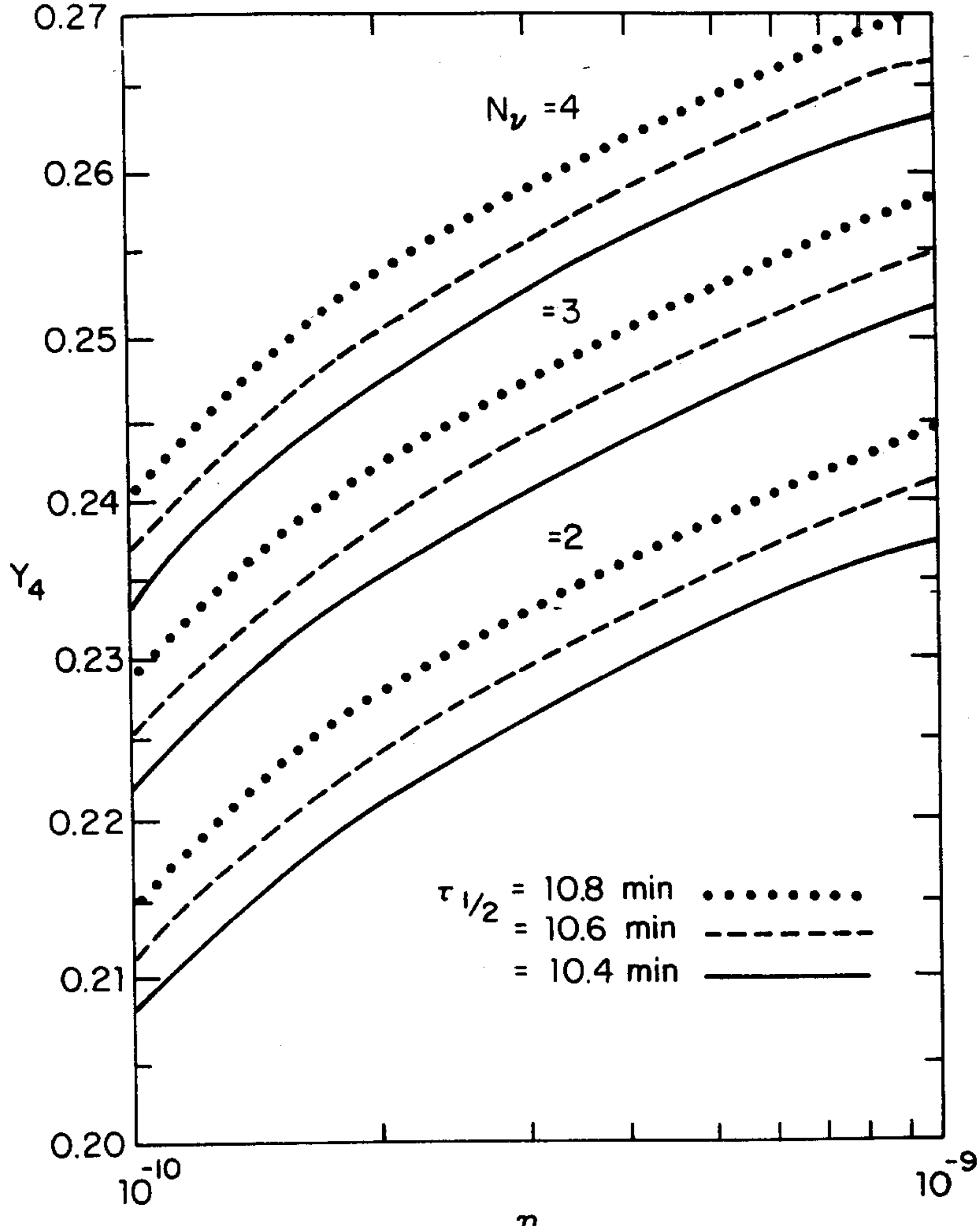
0.27
0.26
0.25
0.24
0.23
0.22
0.21
0.20
Y4
N_ν =4
=3
=2
τ 1/2 = 10.8 min
= 10.6 min
= 10.4 min
10^-10
10^-9
η

FIGURE 6