Vol. B20 (1989) ACTA PHYSICA POLONICA No 5

THREE LECTURES ON PARTICLE PHYSICS AND COSMOLOGY*

By K. A. Olive

School of Physics and Astronomy, University of Minnesota, Minneapolis, MN 55455, USA

(Received October 3, 1988)

In these lectures the particle physics-cosmology connection is reviewed. In the first lecture the standard big bang model, including inflation and baryosynthesis, is outlined. In the second lecture dark matter and some prospects for its detection are discussed. The third lecture explores the role of flat directions in cosmology.

PACS numbers: 98.80.Cq

1. Introduction

The study of the very early Universe in the standard big bang model necessarily requires an intimate connection to particle physics. Already at a late time of the order of a minute, nuclear and particle physics play a crucial role in determining the abundances of the light elements. In fact, agreement with observations requires at most four neutrino species, a limit which is only beginning to be challenged by accelerator experiments. To understand the present baryon asymmetry and the present isotropy seems to require much more than is available by experimental verification. We can only guess at unification models with or without supersymmetry or supergravity or a string theory, and explore their phenomenological and cosmological consequences.

Present-day cosmology also offers a strong connection to particle physics. Through not too well known, the overall mass density of the Universe places constraints on the masses of particles such as neutrinos and photinos. The apparent lack of luminous matter implies that perhaps most of the mass of the Universe is in the form of dark matter. Particle physics provides us with an abundant list of candidates.

In these lectures, I will review the particle physics cosmology connection and concentrate on a set of particle physics models described by scalar potentials with flat directions and examine their cosmological consequences, which include effects on the present mass density, the baryon asymmetry and inflation. In the first lecture I will outline the standard

* Presented at the XXVIII Cracow School of Theoretical Physics, Zakopane, Poland, May 31 — June 10, 1988.

big bang model including inflation and baryosynthesis. In the second lecture I will discuss dark matter and some prospects for detection. In the third lecture I will explore the role of flat directions in cosmology.

LECTURE I

The standard big bang model assumes homogeneity and isotropy and is described by the Friedmann-Robertson-Walker metric [1]

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

where $R(t)$ is the F-R-W scale factor and k is the curvature constant describing the overall geometry ($k = 0, +1, -1$ for a flat, closed or open Universe). Einstein's equations lead to the Friedmann equation which relates the expansion rate to the energy density and curvature,

$$H^2 \equiv (\dot{R}/R)^2 = 8\pi\varrho/3M_P^2 - k/R^2 + \Lambda/3, \tag{2}$$

where H is the Hubble parameter, ϱ is the total mass-energy density, Λ is the cosmological constant and $M_P = G_N^{-1/2} = 1.2 \times 10^{19}$ GeV is the Planck mass.

When combined with the equation for energy conservation (equivalent to entropy conservation)

$$\dot{\varrho} = -3\left(\frac{\dot{R}}{R}\right)(\varrho + p), \tag{3}$$

where p is the isotropic pressure, the Friedmann equation (2) leads to several typical expansion stages in the early Universe. At early times, the Universe is thought to have been dominated by radiation so that the equation of state was just $p = \varrho/3$ and at early times we can assume that the contribution to H from k and Λ were negligible, we have that $\dot{\varrho} = -4\varrho\dot{R}/R$ or $\varrho \sim R^{-4}$ and hence $\dot{R} \sim R^{-1}$ so that

$$R \sim t^{1/2} \tag{4}$$

for a radiation dominated Universe. One then also finds the time temperature relation through

$$t = (3/32\pi G_N\varrho)^{1/2} + \text{constant} \tag{5}$$

and realizing that $\varrho \sim T^4$. Similarly for a matter dominated Universe (late times) we take $p = k = \Lambda = 0$ and find $\dot{\varrho} = -3\varrho\dot{R}/R$ or $\varrho \sim R^{-3}$ and

$$R \sim t^{2/3}. \tag{6}$$

Because $\varrho_{\text{matter}} \sim 1/R^3$ and $\varrho_{\text{rad}} \sim 1/R^4$, at early times $\varrho_{\text{rad}} > \varrho_{\text{matter}}$ and the Universe became matter dominated when T_E ($\varrho_{\text{rad}} = \varrho_{\text{matter}}$) $\simeq 10^3$ K.

If we maintain that $\Lambda = 0$, we can define a critical energy density ϱ_c such that $\varrho = \varrho_c$ for $k = 0$

$$\varrho_c = \frac{3H^2}{8\pi G_N}. \tag{7}$$

In terms of the present value of the Hubble parameter

$$\varrho_c = 1.88 \times 10^{-29} h_0^2 \text{ g cm}^{-3}, \tag{8}$$

where

$$h_0 = H_0/(100 \text{ km Mpc}^{-1} \text{ s}^{-1}) \tag{9}$$

is the present value of the Hubble parameter in units of 100 km Mpc^{-1} s^{-1}. The cosmological density parameter is then defined as the ratio of the present energy density to the critical density

$$\Omega = \varrho/\varrho_c. \tag{10}$$

Furthermore, the value of Ω will determine the sign of k. For $\Omega > 1$ we have $k = +1$, $\Omega = 1$ corresponds to $k = 0$ and $\Omega < 1$ to $k = -1$. In terms of Ω the Friedmann equation can be rewritten as

$$(\Omega - 1)H^2 = \frac{k}{R^2}. \tag{11}$$

The observational limits on h_0 and Ω are [2].

$$0.4 < h_0 < 1, \tag{12a}$$

$$0.1 < \Omega < 4. \tag{12b}$$

One of the most important pieces of evidence we have for the big bang model is the existence of the 3 K microwave background radiation [3]. Since the formation of neutral hydrogen (recombination) the observed photons have been decoupled. They are a remnant of a previous hot and thermal era. The energy density in this background is simply $\varrho_\gamma = (\pi^2/15)T^4$ (in units where $\hbar = c = k_B = 1$). Recombination occurs when $T \sim 1$ eV or close to T_E. At higher temperatures, the Universe is radiation dominated and the total energy density is given by

$$\varrho = \left(\sum_B g_B + \tfrac{7}{8} \sum_F g_F\right) \frac{\pi^2}{30} T^4 = \frac{\pi^2}{30} N(T) T^4, \tag{13}$$

where $g_{B(F)}$ are the total number of relativistic boson (fermion) degrees of freedom. For temperatures $T \gtrsim 10^2$ GeV, $N(T)$ becomes very model dependent.

A second piece of evidence in support of the big bang model is the agreement between observations of the abundances of the light elements and the predictions of big bang nucleosynthesis [4–6]. At very high temperatures ($T \gg 1$ MeV), the weak interaction rates for the processes

$$n+\nu_e \leftrightarrow p+e^-,$$
$$n+e^+ \leftrightarrow p+\nu_e,$$
$$n \leftrightarrow p+e^- +\bar{\nu}_e, \tag{14}$$

were all in equilibrium, i.e., $\Gamma_{wk} \sim H$. Thus we would expect that initially (n/p) = 1. Actually in equilibrium, the ratio is essentially controlled by the Boltzmann factor so that

$$(n/p) = \exp(-\Delta m/T), \tag{15}$$

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; $\bar{E}_\gamma = 2.7$ T 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

$$\eta^{-1} \exp[-2.2 \text{ MeV}/T], \tag{16}$$

where $\eta = n_B/n_\gamma$. When this quantity (16) becomes $\lesssim O(1)$, the rate for $p+n \to D+\gamma$ finally becomes greater than the rate for dissociation $D+\gamma \to p+n$. This occurs when $T \sim 0.1$ MeV or when the Universe is a little over 2 min. old.

Because the rates for processes (14) freeze out at $T \sim 1$ MeV, the neutron to proton ratio must be adjusted from its equilibrium value at freeze out to its value when nucleosynthesis begins. When freeze out occurs, the ratio (n/p) is relatively fixed at

$$(n/p) \sim 1/6. \tag{17}$$

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

$$(n/p) \sim 1/7. \tag{18}$$

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

$$X_4 \equiv (N_{^4He}/N_H) = 1/2\,(n/p)/[1-(n/p)] \tag{19}$$

or more importantly the ^{4}He mass fraction

$$Y_4 = 4X_4/(1+4X_4) = 2(n/p)/[1+(n/p)]. \tag{20}$$

For $(n/p) \simeq 1/7$, we estimate that $Y_4 = 0.25$ 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 [5, 6]. 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.

There are three 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 (16). Basically what happens is that for a larger baryon-to-photon ratio η the quantity (16) 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. (20)] 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.

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

$$n_B = \varrho_B/m_B = \Omega_B\varrho_c/m_B = 1.13\times 10^{-5}\Omega_B h_0^2\ \text{cm}^{-3}, \tag{21}$$

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

$$n_\gamma = 400\,(T_0/2.7)^3\ \text{cm}^{-3}, \tag{22}$$

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

$$\eta = 2.81\times 10^{-8}\Omega_B h_0^2(2.7/T_0)^3. \tag{23}$$

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 best value of $\tau_{1/2} = 10.4$ min. is actually uncertain by about two percent and this is enough to affect the production ^{4}He. $\tau_{1/2}$ can be as low as 10.2 min. As in the case of η, increasing $\tau_{1/2}$ leads to a larger value of Y_4.

The final input parameter, I said was the number of light (relativistic), stable ($\tau \gtrsim 1$s) particles. Likely candidates for these particles are neutrinos and thus the number of neutrino flavors N_ν becomes important. Of course any other type of light particles such as majorans

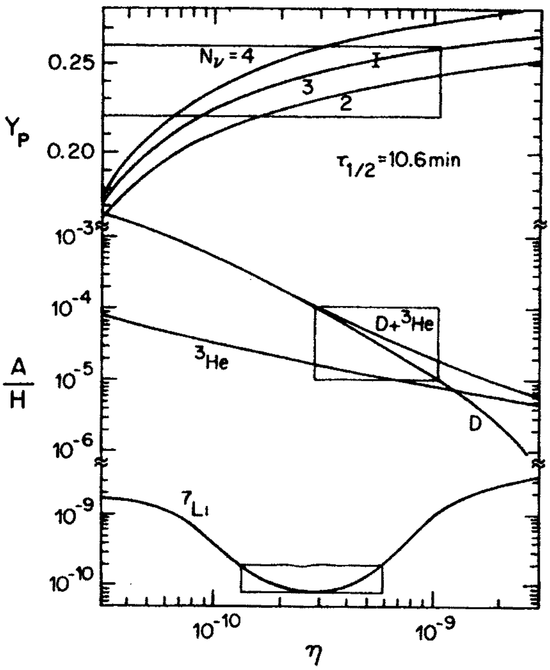

Fig. 1. The predicted primordial abundances of ^{4}He (by mass), D, ^{3}He and ^{7}Li (by number relative to H) as a function of η for $\tau_{1/2}$ = 10.6 min; for ^{4}He the predictions for N_ν = 2, 3, 4 are shown and the size of the "error" bar shows the range in Y_p which corresponds to $10.4 < \tau_{1/2} < 10.8$ min. The boxes show the ranges in abundance determined by observations yielding limits on η

or axions, etc., may also be important. The number of neutrino flavours N_ν will also affect the primordial abundance of ^{4}He and like η and $\tau_{1/2}$, increasing N_ν increases Y_4.

In Fig. 1, the predicted abundances of D, ^{3}He, ^{4}He and ^{7}Li are plotted as a function of η. Also shown in the figure [6] are the observational ranges for these abundances. From deuterium alone, D/H $\gtrsim$ 1–2 $\times 10^{-5}$ we see that $\eta \lesssim$ 7–10 $\times 10^{-10}$ while the combination of (D + ^{3}He)/H $\lesssim 10^{-4}$ indicates that $3 \times 10^{-10} \lesssim \eta$. The constraint from ^{7}Li is consistent (^{7}Li/H) $\lesssim 2 \times 10^{-10}$ implies $10^{-10} \lesssim \eta \lesssim 7 \times 10^{-10}$. Most importantly the range $0.22 \lesssim Y_4 \lesssim 0.25$ or at most $Y_4 \lesssim 0.26$ is again consistent with this same range for η.

Figure 1 actually contains significantly more information than just a limit on η. We can set a limit [6–8] on N_ν provided that we have a lower limit to η. Using $\eta \geqslant 3 \times 10^{-10}$ and $Y_4 \leqslant 0.25$ (0.26), we find that $N_\nu \leqslant 4$ (4.6) 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, $\eta \geqslant 3 \times 10^{-10}$ and $\tau > 10.4$ min. allows $N_\nu = 4$ only if $Y_4 \geqslant 0.253$; however, we are not yet in a position to believe the third decimal place. We can also turn the limits around and set a lower limit to the helium abundance by assuming $\eta \geqslant 3 \times 10^{-10}$ and $N_\nu \geqslant 3$ 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/or for unstable (the present limit is only $m_{\nu_\tau} < 70$ MeV). If we only assume $N_\nu \geqslant 2$, then the lower limit on Y_4 becomes $Y_4 \geqslant 0.22$. 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

$$3\times 10^{-10} \leqslant \eta \leqslant 7\times 10^{-10} \tag{24}$$

can be converted to a limit on the baryon density and Ω_B. If we turn around Eq. (23) we have

$$\Omega_B = 3.56\times 10^7 \eta h_0^{-2}(T_0/2.7)^3, \tag{25}$$

and using the limits on η, Eq. (24), h_0, Eq. (12a), and T_0 from (2.7–2.8) K we find a range for Ω_B

$$0.01 \leqslant \Omega_B \leqslant 0.17. \tag{26}$$

Recall that for a closed Universe $\Omega > 1$, thus from Eq. (26) we can conclude that the Universe is not closed by baryons. Indeed if $\Omega_B = 1$ and $h_0 > 0.4$, $T_0 < 2.8$ K we have $\eta > 4\times 10^{-9}$ and the following light element abundances: D/H $\lesssim 10^{-7}$, ^{3}He/H $\lesssim 5\times 10^{-6}$, ^{7}Li/H $\gtrsim 5\times 10^{-9}$ and $Y_4 \gtrsim 0.269$. 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 $\Omega = 1$ as implied by inflation the limit from nucleosynthesis would indicate that some form of dark matter must exist.

Recently, the possibility that fluctuations in the matter density produced in the quark--hadron confinement transition [9, 10] could affect the primordial yields has been explored [11]. The idea is that because of different diffusion lengths for neutrons and protons, the neutron density could smooth itself out leaving inhomogeneities in the proton density. Nucleosynthesis would begin with a non-uniform n/p ratio. It was hoped that perhaps $\Omega_B = 1$ would yield abundances in agreement with observations. The biggest problem first was the ^{7}Li abundance which was 1–2 orders of magnitude too large. It was hoped then that nucleosynthesis might be able to place constraints [12, 13] on the magnitude of the baryon fluctuations, $(n_B)_Q/(n_B)_H \lesssim 4$. However estimates of these fluctuations found that $(n_B)_Q/(n_B)_H \gtrsim 6$–7 for all transitions temperatures $T_c \sim 100$–1000 MeV. More recently, it has become apparent that diffusion during nucleosynthesis [14] also plays an important role and implying that for $\Omega_B = 1$ all abundances might be in conflict with observation. It seems at the present time that the predictions of standard nucleosynthesis are safe.

The origins of the current connection between particle physics and cosmology really began with the generation [15] of a small but finite baryon to entropy ratio using grand unified theories (GUTs). 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

$$\eta = (n_B - n_{\bar{B}})/n_\gamma \tag{27}$$

(if $n_{\bar{B}} = 0$, this η is the same as that used in the previous discussion) since antimatter is not observed in primary form,

$$\eta \sim n_B/n_\gamma \sim 10^{-10}-10^{-9}. \tag{28}$$

In a standard model, the entropy density today is related to n_γ by

$$s \simeq 7n_\gamma \tag{29}$$

so that Eq. (28) implies $n_B/s \sim 10^{-10}-10^{-11}$. This ratio is conserved however and hence represents an 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 of $B\bar{B}$ 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 $\Gamma_A \simeq n_N m_\pi^{-2}$ fell below the expansion rate. This occurred at $T \simeq 20$ MeV. However, at this time the nucleon number density has already dropped to

$$n_N/n_\gamma = n_{\bar{N}}/n_\gamma \simeq 10^{-18} \tag{30}$$

which is eight orders of magnitude too small [16] 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 $T = 20$ MeV, the age of the Universe was only $t = 2 \times 10^{-3}$ 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, $10^{12}\ M_\odot$.

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 nonzero value of η using GUTs.

There are three basic ingredients necessary [15] 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 anti-baryons? 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 $X \to e^- + d$, $\bar{u} + \bar{u}$. Although today these interactions are very ineffective because of the masses of the X and Y bosons, in the early Universe when $T \sim M_X \sim 10^{15}$ 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 ane 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 [17]. 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

$$\Gamma_D \sim \alpha M_X. \tag{31}$$

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

$$\alpha M_X \gtrsim N(T)^{1/2} T^2 / M_P \tag{32}$$

or at a temperature

$$T^2 \lesssim \alpha M_X M_P N(T)^{-1/2}. \tag{33}$$

Scatterings on the other hand proceed at a rate

$$\Gamma_S \sim \alpha^2 T^3 / M_X^2 \tag{34}$$

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

$\overline{X}$'s). Thus the condition for equilibrium is that at $T = M_X$, $\Gamma_D > H$ or

$$M_X \lesssim \alpha M_P (N(M_X))^{-1/2} \sim 10^{18} \alpha \text{ GeV}. \quad (35)$$

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 $n_X/n_\gamma \sim (M_X/T)^{3/2} \exp[-M_X/T] \ll 1$.

Hence, the decays are occurring out of equilibrium (inverse decays are not occurring), and we have the possibility for producing a net asymmetry.

If total baryon density that will have been produced by the X, $\overline{X}$ pair [provided Eq. (35) is not satisfied] is

$$n_B \sim (\Delta B) n_X, \quad (36)$$

where ΔB is the baryon number produced by a X, $\overline{X}$ pair, then since we also have $n_X = n_{\overline{X}} \sim n_\gamma$,

$$n_B \simeq (\Delta B) n_\gamma. \quad (37)$$

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

$$s = \frac{2\pi^2}{45} N(T) T^3. \quad (38)$$

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

$$n_B/s \sim 10^{-2} (\Delta B). \quad (39)$$

The value of n_B/s that we are looking for must be related to the limits on η from nucleosynthesis. η in the range $(3\text{–}7) \times 10^{-10}$ corresponds to a value of n_B/s in the range $(4.3\text{–}10) \times 10^{-1}$. Comparing this with the expected production, Eq. (39) 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.

In Figs. 2a and 2b, results [18] for the development and final baryon asymmetries are shown for an SU(5) model. An alternative scenario for baryosynthesis will be discussed in Lecture III.

Closely tied to particle physics is cosmological inflation. Inflation as will be described below is the effect of a phase transition needed to solve a host of cosmological problems. 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

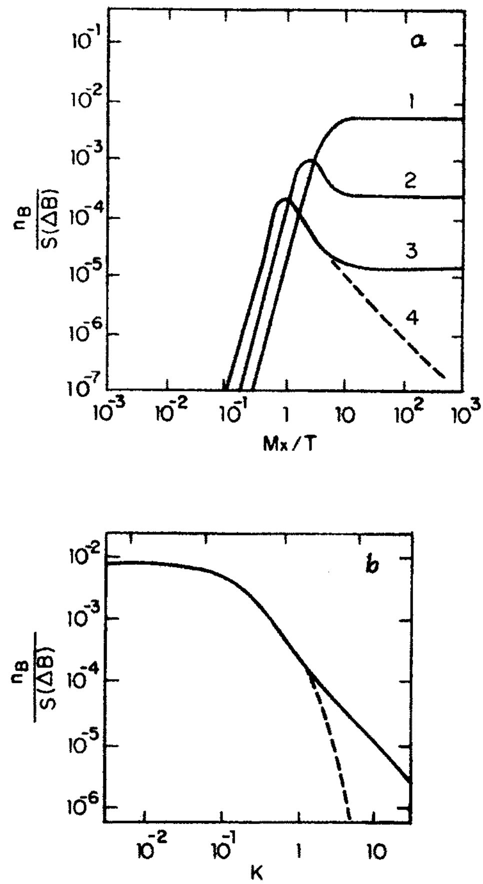

Fig. 2. a) The time evolution of the baryon asymmetry inu nits of (ΔB) for *1* — $M_x = 3\times10^{18}\alpha$; *2* — $M_x = 3\times10^{17}\alpha$; *3* — $M_x = 3\times10^{16}\alpha$; *4* — if scatterings remain very effective; b) The final baryon asymmetry as a function of $K = 3\times10^{17}\alpha/M_x$ in units of (ΔB). The dashed curve assumes effective scatterings

with the 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

$$V_0'/V_d \sim (V_0/V_d)\,(T_0/T_d)^3 \sim (t_0/t_d)^3\,(T_0/T_d)^3 \sim 10^5, \tag{40}$$

where I have used $t_d \sim 3\times10^{12}$ sec and $t_0 \sim 5\times10^{17}$ sec. The ratio (40) corresponds to the number of regions that were casually disconnected at recombination which grew into our present visible Universe. Because the anisotropy of the microwave background is so small, $\delta T/T <$ few $\times\ 10^{-5}$, 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 (2) for the expansion of the Universe and use the limits $\Omega < 4$ and $H_0 < 100\,\mathrm{km\,s^{-1}\,Mpc^{-1}}$ we can form a dimensionless constant

$$\hat{k} = k/R^2T^2 = (\Omega-1)H_0^2/T_0^2 < 3H_0^2/T_0^2 < 2\times 10^{-58}, \tag{41}$$

where I have used $T_0 \geqslant 2.7$ K. In an adiabatically expanding Universe, $\hat{k}$ is absolutely constant ($R \sim T^{-1}$) and thus the limit (41) 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 $T \sim 10^{-1}\,M_{\mathrm{P}}$. 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 (41) 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 among the two main problems that led Guth [19] 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 scalar potential describing a phase transition from a symmetric false vacuum state $\langle\Sigma\rangle = 0$ for some scalar field Σ to the broken true vacuum at $\langle\Sigma\rangle = v$ as in Fig. 3, 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

$$(RT)_f \sim (T_R/T_c)\,(RT)_i \tag{42}$$

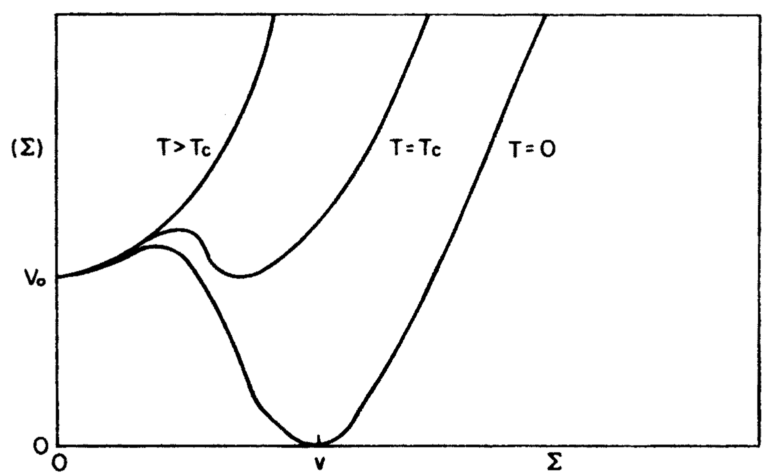

Fig. 3. Schematic view of the scalar potential for a first order phase transition. For SU(5) symmetry breaking Σ is the adjoint and and $V \simeq M_x$

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

The inflationary Universe scenario [19], is based on just such a situation. If during some phase transition, such as $SU(5) \to SU(3) \times SU(2) \times U(1)$ (taking Σ to be in the adjoint representation of SU(5)) the value of RT changed by a factor of $O(10^{29})$, 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 $O(1)$ and have been driven small by the expansion.

If, in an extreme case, a barrier as in Fig. 3 caused a lot of supercooling such that $T_c^4 \ll V(0)$, the dynamics of the expansion would have greatly changed. In the example of Fig. 3 the energy density of the symmetric vacuum, $V(0)$ acts as a cosmological constant with

$$\Lambda = 8\pi V(0) M_p^2. \tag{43}$$

If the Universe is trapped inside the false vacuum with $\langle \Sigma \rangle = 0$, eventually the energy density due, to say, radiation will fall below the vacuum energy density, $\varrho \ll V(0)$. When this happens, the expansion rate will be dominated by the constant $V(0)$ and we will get the De Sitter-type expansion and from Eq. (2)

$$R \sim \exp [Ht], \tag{44}$$

where

$$H^2 = \Lambda/3 = 8\pi V(0)/3M_P^2. \tag{45}$$

The cosmological problems could be solved if

$$H\tau > 65, \tag{46}$$

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

$$\frac{\pi^2}{30} N(T_R) T_R^4 = V(0), \tag{47}$$

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

In the original inflationary scenario, the phase transition given by a potential with a large barrier as in Fig. 3 proceeds via the formation of bubbles [20]. The Universe hopefully reheats, and the release of entropy must occur through bubble collisions and the transition is completed when the bubbles fill up all of space. It is now known [21], however, that the requirement for a long timescale τ is not compatible with the completion of the phase transition. The Universe as a whole remains trapped in the exponentially expanding phase containing only a few isolated bubbles of the broken SU(3) × SU(2) × U(1) phase.

The well-known solution to this dilemma is called the new inflationary scenario [22]. New inflation is based on symmetry breaking using a flat potential of type shown in Fig. 4. In addition to producing $\Omega = 1$, new inflation is capable of producing scale invariant density perturbations [23] of the type preferred for galaxy formation models. However, the original [22] new inflationary models based on a Coleman-Weinberg [24] type of SU(5) breaking produced density fluctuations with magnitude $\delta\varrho/\varrho \sim 50$ rather than $\delta\varrho/\varrho \sim 10^{-4}$ as needed to remain consistent with microwave background anisotropies.

This problem led to primordial inflation [25] and the inception of the inflaton [26] as a new scalar field whose sole (hopefully temporary) role is to drive inflation. Primordial

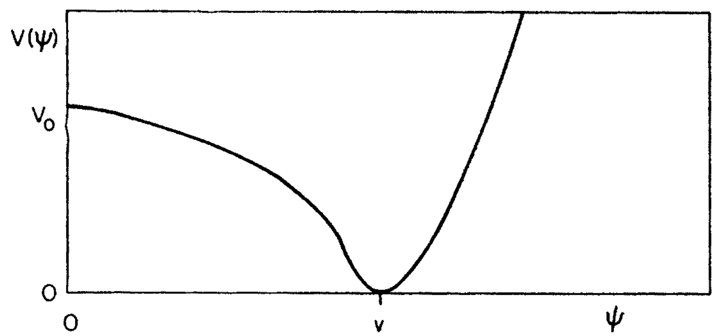

Fig. 4. Schematic view of the scalar potential for new inflation. In primordial inflation, ψ, is the inflaton and $v \sim M_P$

inflation, allowed the expectation value v to go up to the Planck scale producing a flatter potential. Supersymmetry was also employed [27, 25] to give flatter potentials and acceptable density perturbations [28]. These models were then placed in the context of $N = 1$ supergravity [29, 30]. Additional problems regarding initial conditions [31–33] were resolved in non-minimal versions [34–36] of primordial inflation [37].

These models have the benefit that they can be written down very simply. We can assume a scalar potential of the form $V(\psi) = m^4 P(\psi)$, where ψ is the inflaton and $P(\psi)$ is some polynomial with coefficients $O(1)$. Such a potential would arise from a superpotential of the form

$$f(\psi) = m^2(1-\psi/M_P)^2 M_P, \tag{48a}$$

$$f(\psi) = m^2(\phi-\psi^4/4M_P^3), \tag{48b}$$

where (48a) [30] is to be used in minimal supergravity while (48b) [45] is to be used in SU(N, 1) supergravity. The former does not satisfy thermal initial conditions [31] while the latter does. During inflation, the Hubble parameter is

$$H^2 \simeq m^4/M_P^2. \tag{49}$$

The magnitude of density perturbations is given by

$$\delta\varrho/\varrho \sim (m^4/M_P^3 H)\ln^2(Hk^{-1}) \sim 10^3 m^2/M_P^2, \tag{50}$$

where k^{-1} is the length scale of the perturbation. $\delta\varrho/\varrho \sim 10^{-4}$ requires $m^2 \sim 10^{-7} M_P^2$ so that the amount of inflation, $\exp(H\tau)$, is determined by

$$H\tau \sim M_P^2/m^2 \sim 10^7. \tag{51}$$

Clearly more than enough inflation.

Other variants of the inflationary scenario include chaotic inflation [38] and stochastic inflation [39]. Unfortunately to date there are no models of inflation in the context of string theory [40].

LECTURE II

If inflation gives us $\Omega = 1$ while nucleosynthesis restricts us to $\Omega_B \lesssim 0.2$, the bulk of the mass density of the Universe is in some form of dark matter (DM). Already on galactic scales, there is good evidence from rotation curves [41] of spiral galaxies for the presence of dark matter and a galactic halo. The rotation curve is a measure of the velocity as a function of distance from the center of the galaxy of a star as it revolves around the galaxy. If there were no DM, one would expect that at distances beyond the bulk of the luminous matter that $v^2 \sim 1/r$. Instead one finds flat rotation curves ($v^2 \sim$ constant) out to very large distances (> 50 kpc). This implies that the mass of the galaxy must continue to increase $M \sim r$ beyond the luminous region.

A more subtle DM problem is the one in relation to the growth of density perturbations and galaxy formation. One of the features of the spectrum of density perturbations produced by inflation is that as the different Fourier modes fall within the horizon scale (i.e., $\lambda \sim ct$) they all have the same magnitude $\frac{\delta\varrho}{\varrho}$ at that time. Once within the horizon, these modes cannot really begin to grow further as long as the Universe is radiation dominated. At a temperature of a few thousand degrees, the Universe becomes matter dominated and density perturbations begin growing as

$$\frac{\delta\varrho}{\varrho} \sim R(t) \sim \frac{1}{T}. \tag{52}$$

Now to reach non-linear growth we must have had $\frac{\delta\varrho}{\varrho} \lesssim 1$ at the time when the oldest galaxies and quasars were forming or at $T \sim 4\,T_0 \sim 10$ K. This means that at the time of matter dominance $\frac{\delta\varrho}{\varrho} \sim 10^{-3}$. However we know from limits on the anisotropy of the microwave background that $\frac{\delta\varrho}{\varrho} < 10^{-4}$. We will see shortly that this is really a DM problem.

As a first guess as to the identity of the DM, one might pick baryons i.e., ordinary matter. As we saw from big bang nucleosynthesis, there are good limits (Eq. (26)) on the value of Ω in the form of baryons. Recall, in the standard model, one finds good agreement for the predicted abundances of the light isotopes D, ^{3}He. ^{4}He, ^{7}Li only for a range in Ω_B between 0.01 and 0.17. For $\Omega_B < 0.01$, D and ^{3}He are overproduced while for $\Omega_B > 0.17$, ^{4}He is overproduced and D is underproduced. If inflation is correct and $\Omega = 1$, then at least some of the DM must be non-baryonic.

Returning to the growth of perturbations, if there exists some form of non-baryonic DM, the Universe may have become matter dominated earlier. For example in the case of massive neutrinos $T_{MD} \sim m_\nu/10$. Density perturbations could then begin to grow earlier at say $T \sim 10^4$–$10^5\ T_0$ while baryonic perturbations could not until decoupling at $T \sim 10^3\ T_0$. After decoupling, the baryons would fall into the perturbation already formed by the neutrinos. Hence the existence of DM could help enormously in the growth of perturbations for galaxy formation. For a complete review see Ref. [42].

Returning once more to the question of DM on the scale of galactic halos, although one needs $\Omega \gtrsim 0.05$ and that is consistent with Ω in baryons, there are several arguments [43] against baryonic matter in halos. Put briefly, it is very difficult to have a large baryon density in such a way that it is unobservable. In the form of gas the baryons would heat up and emit X-rays in violation of observed limits. To put the baryons in non-nuclear burning stars (Jupiters) would require an extrapolation of the stellar mass distribution which is very different from what is observed. Dust or rocks along with dead remnants such as neutron stars or black holes would require a metal abundance in great excess of the galactic metallicity. Very massive ($\gtrsim 100\ M_\odot$) black holes remain a possibility.

There are of course, many other candidates for the DM. Because of its important role in the formation of galaxies, DM has classified [44] into three types: hot, warm and cold DM. They are distinguished by their effective temperature at the time they decoupled from the thermal background. Examples of hot particles are neutrinos or very light Higgsinos with < 100 eV masses. These particles decouple at $T_d \sim 1$ MeV and are thus still relativistic at T_d. They typically produce large scale structure first. Warm particles decouple earlier and have higher masses (up to ~ 1 keV). Any superweakly interacting neutral particle is a warm candidate such as a right-handed neutrino. Cold particles are non-relativistic at temperatures relevant for galaxy formation and usually have masses > 1 GeV. Examples of these include heavy neutrinos, photinos/Higgsinos, sneutrinos and axions. These typically produce small scale structure first. With this classification, the specific identity of the particle is no longer important for the purposes of galaxy formation.

Given the need for DM, we can ask what sort of constraints are there on particle properties. The most common cosmological constraint is on the mass of a stable particle and is derived from the overall mass density of the Universe. The mass density of a particle x can be expressed as

$$\varrho_x = m_x Y_x n_\gamma \leqslant \varrho_c \simeq 10^{-5} h_0^2\ \mathrm{GeV/cm^3}, \tag{53}$$

where $Y_x = n_x/n_\gamma$ is the density of x's relative to the density of photons, for $\Omega h_0^2 < 1$. Hot particles have limits characteristic to that of neutrinos. For neutrinos [45] $Y_\nu = 3/11$ and one finds

$$\sum \left(\frac{g_\nu}{2}\right) m_\nu < 100\,\text{eV}\,(\Omega h_0^2), \tag{54}$$

where the sum runs over neutrino flavors and $g_\nu = 2$ for a Majorana mass neutrinos and $g_\nu = 4$ for a Dirac mass neutrino. All hot particles with abundances Y similar to neutrinos will have mass limits as in Eq. (54).

Warm particle limits are derived from Eq. (54) as well. Warm particles have lower abundances than neutrinos and the corresponding mass limits are weaker. Recall that $Y_\nu = 3/11$ is derived from the conservation of entropy before and after $e^\pm$ annihilation. Neutrinos at this time are decoupled so that after the annihilations $(T_\nu/T_\gamma)^3 = 4/11$ and $Y_\nu = (3/4)\,(T_\nu/T_\gamma)^3 = 3/11$. (The factor of 3/4 is due to the difference between Fermi and Bose statistics.) If a particle x interacts more weakly than neutrinos then the ratio $(T_x/T_\gamma)^3$ will be lowered [46] due to other particle species annihilations. Thus Y_ν is reduced allowing [47] for a larger value for m_x. If the particle x decouples around the GUT epoch, then Y_x could be as low as $O(10^{-2})$ and $m_x < O(1)$ keV.

For cold particles the analysis is somewhat different. The abundance is now a function of m_x and in most cases one finds a lower limit to m_x. The reason for this is that for large m_x, Y_x is controlled by the annihilations of x. When the annihilations freezeout, Y_x is fixed. The freezeout will then depend on the annihilations cross-section and roughly one finds $Y_x \sim (m_x\sigma_A)^{-1}$ and $\varrho_x \sim 1/\sigma_A$. This situation was first analyzed for neutrinos [48]. The annihilation cross-section in this case is basically $\sigma_A \sim m_\nu^2/m_W^4$ so that $\varrho_x \sim 1/m_\nu^2$ and yields [48–50] $m_\nu > 4$ GeV for Dirac mass neutrinos and $m_\nu > 6$ GeV for Majorana mass neutrinos [51, 49, 50].

Supersymmetric theories introduce several DM candidates. The reason is that if the R-parity (which distinguishes between "normal" matter and the supersymmetric partners) is unbroken then there is at least one supersymmetric particle which must be stable. Candidates for the stable particle include the photino, Higgsino, and sneutrino. If we assume for simplicity that all of the scalar quarks and leptons have equal masses then the photino annihilation cross-section can be expressed as [52–54]

$$\langle \sigma v \rangle_A \simeq \frac{8\pi\alpha^2}{m_{sf}^4} \sum q_f^4 (1 - z_f^2) m_{\tilde\gamma}^2 (z_f^2 + 2x(1 - 5z_f^2/2)), \tag{55}$$

where α is the fine structure constant, m_{sf} is scalar fermion mass, q_f the electric charge of the fermion f, $z_f = m_f/m_{\tilde\gamma}$ and $x = T/m_{\tilde\gamma}$. For $m_{sf} \sim 90$ GeV, $m_{\tilde\gamma} > 12$ GeV. For Higgsinos [54] the annihilations are controlled by the fermion Yukawa couplings and the cosmological bound requires $m_{\tilde H} > m_b$ or about 5 GeV.

Sneutrinos are an interesting example in that there is in general no cosmological limit on their mass [55]. In addition to the standard weak annihilations of sneutrinos, there is also the process $\tilde\nu + \bar{\tilde\nu} \to \nu + \bar\nu$ via zino exchange. In this case $\langle\sigma v\rangle_A \propto 1/M_{\tilde Z}$ and is independent of m_ν. Thus ϱ_ν is fixed by parameters other than m_ν making the sneutrino mass free from cosmological bounds.

Assuming the existence of dark matter and assuming that the dark matter is responsible for the observed flat rotation curves of galaxies it is possible to estimate the mass density of dark matter in an isothermal halo,

$$\varrho = \frac{v_r^2}{4\pi G}\frac{1}{a^2+r^2} = 0.07\ \mathrm{GeV\ cm^{-3}}\frac{v_{100}^2}{a_{10}^2+r_{10}^2}, \tag{56}$$

where $v_r = 100\, v_{100}$ km s^{-1} is the galactic rotation velocity, $a = 10a_{10}$ kpc is some core radius for the halo and $r = 10\, r_{10}$ kpc is the distance from the galactic center to the sun. Typical values are $v_{100} \simeq 2.4$, $r_{10} \simeq 1$ and $a_{10} \simeq 0.6$ so that $\varrho = 0.3$ GeV cm^{-3}.

Dark matter particles in the solar neighborhood may be trapped [56] in the sun as they pass through and elastically scatter. Numerically, the trapping rate is [56, 57].

$$\Gamma_t \simeq 10^{29}\ \mathrm{s^{-1}}(n_x/0.3\ \mathrm{cm^{-3}})(300\ \mathrm{km\ s^{-1}}/\bar{v})(1\ \mathrm{GeV}/m_x)(\sigma_E/10^{-36}\ \mathrm{cm^2}), \tag{57}$$

where n_x is the dark matter abundance in the solar neighborhood, $\bar{v}$ is the R.M.S. velocity of x in the halo and σ_E is the elastic scattering cross-section.

The abundance of dark matter particles in the sun is controlled by annihilations [58, 59] and evaporation [60, 56, 61] which is negligible for $m_x \gtrsim 3$ GeV. Although the annihilations make it difficult [59] to resolve to the solar neutrino problem by heat transfer [60] they open up a possibility for detecting a signature of the presence [58, 57, 62–67] of dark matter. The products of these annihilations are high energy neutrinos. In particular, we will be interested in looking for the prompt neutrinos in reactions such as [62, 63]

$$\begin{aligned} \mathrm{XX} &\to \mathrm{f\bar{f}} \\ &\to \mathrm{f'} + \bar{\mathrm{l}} + \nu_l . \end{aligned} \tag{58}$$

It is then straightforward to calculate the differential flux of neutrinos produced in this way [62]

$$\frac{d\Phi}{dE_\nu} = \sum_f \frac{\frac{1}{2}\Gamma_t B_f}{4\pi d^2}\frac{1}{\Gamma_f}\frac{d\Gamma_f}{dE_\nu}, \tag{59}$$

where $d = 1$ A.U., B_f is the branching ratio for $\mathrm{f} \to \mathrm{f'} + \bar{\mathrm{l}} + \nu_l$ and Γ_f is the decay rate for the same process. We can now compare Eq. (59) with the differential flux of atmospheric neutrinos produced by cosmic rays [68]. In Figs. 5a–d [62], the differential neutrino flux for a) (generic) higgsinos (these are equivalent to Majorana mass neutrinos, i.e. their abundance is completely determined by their mass); b) (closure) higgsinos (their abundance can always be adjusted so that $\Omega = 1$); c) photinos; and d) also shown is a typical candidate from superstring inspired theories [69]. Indicated on the figures, is the mass of each particle in the range 6–40 GeV (a–c) and 20–50 GeV (d). The dashed line is the atmospheric background flux of ν_e's within 30° of the sun. Also shown in Fig. 6 is the total flux of monochromatic neutrinos [57] for the Dirac mass neutrinos and sneutrinos. The differential flux

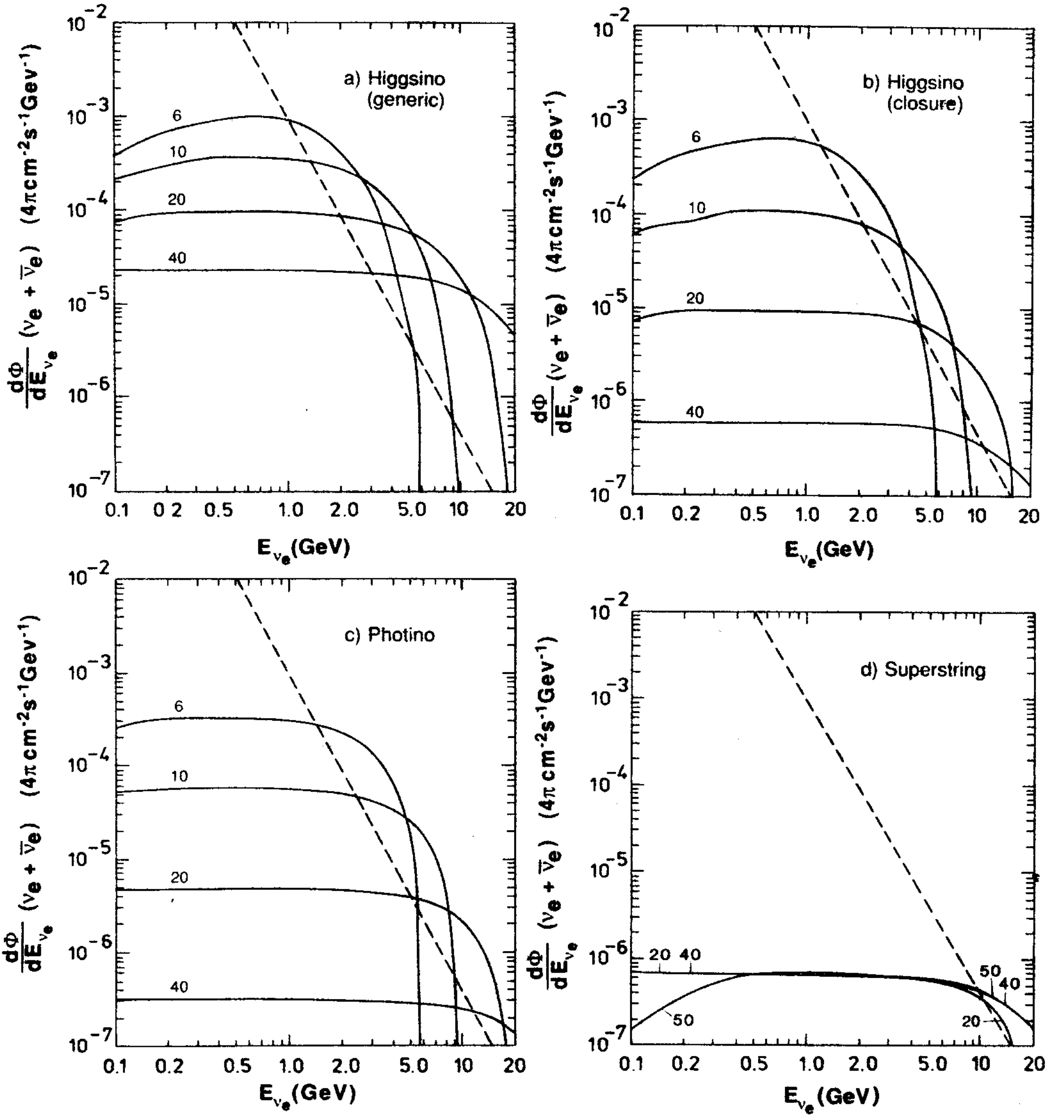

Fig. 5. The differential flux of $\nu_e + \bar{\nu}_e$ from the sun due to the annihilation of cold dark matter (χ) as compared to the atmospheric background flux in the direction of the sun (dashed line) for: a) a generic higgsino, $\chi = \tilde{H}$ or $\tilde{\bar{H}}$ for $m_\chi = 6$, 10, 20 and 40 GeV. In this case, $\Omega_\chi = 1$, 0.37, 0.12 and 0.36 respectively; b) a symmetric higgsino, $\chi = (\bar{v}\tilde{H} + v\tilde{\bar{H}})/(v^2 + \bar{v}^2)^{1/2}$ with $v/\bar{v}$ adjusted so that $\Omega_\chi = 1$: $v/\bar{v} = 3$, 1.8, 1.4 and 1.2 for the same set of LSP masses as in a); c) a photino with degenerate squark and slepton masses, $m_{\tilde{f}}$, adjusted so that $\Omega_\chi = 1$: $m_{\tilde{f}} = 71$, 88, 117 and 166 GeV for the same set of LSP masses as in a); d) superstring LSP (photino) with masses 20, 40 and 50 GeV and sfermion masses taken from the relations of Refs. [69, 101]

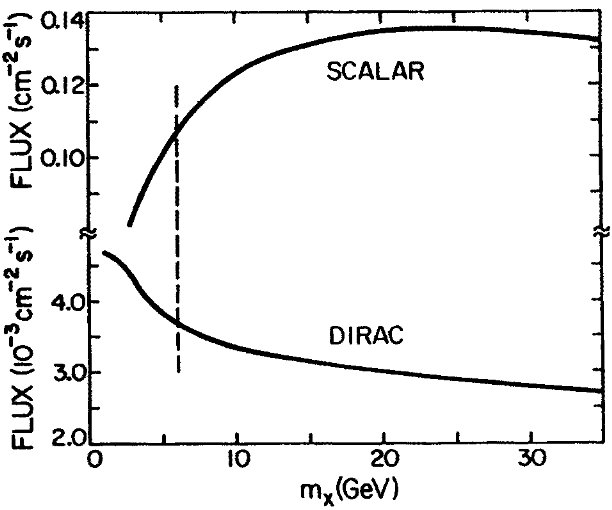

Fig. 6. The total flux of monochromatic $\nu_e + \bar{\nu}_e$ produced by annihilation in the sun of scalar or Dirac neutrinos

for these particles is just

$$\frac{d\Phi}{dE} = \Phi\delta(E - m_x). \tag{60}$$

In order to make the comparison between the calculated dark matter induced flux and the background we must compute an event rate

$$\bar{S} \propto \int dE\, V_F(E)\sigma_\nu(E)\, \frac{d\Phi}{dE_\nu} \tag{61}$$

for neutrinos from the sun, where V_F is the fiducial volume of the detector and σ_ν is the neutrino cross-section in the detector. A similar computation is done for the atmospheric background rate $\bar{A}$. We also distinguish between three types of events [65]: 1) contained events with 1 GeV $\leqslant E_\nu \leqslant$ 2 GeV; 2) neutrino events with $E_\nu \geqslant$ 2 GeV producing through--going muons; and 3) contained events with $E_\nu \geqslant$ 2 GeV.

Data for each of these event types exists. For case 1), IMB [70] reports 11 events within 30° of the sun out of a total of 89 events. The 90% confidence level statistical upper limit on the ratio $r \equiv \bar{S}/\bar{A}$ is $r \leqslant 0.14$ [64, 65]. For case 2), IMB [70] reports 2 through-going muon events within 8° of the sun out of a total of 187 events yielding $r \leqslant 0.024$ at the 90% confidence level. Finally for case 3), IMB [70] reports 0 events out of 10 within 30° of the sun for $r \leqslant 0.23$; Kamioka [71] reports 0 events out of 23 for $r \leqslant 0.10$ and Frejus [72] reports 0 events out of 24 with $r \leqslant 0.096$. Because the data is still statistics limited, the best limit for contained events with $E_\nu \geqslant$ 2 GeV comes from combining the data yielding [65]

$$r < 0.040. \tag{62}$$

In the tables, the calculated values of r are shown [65].

Before discussing the limits we can set based on this data, it is worthwhile to remind the reader of the uncertainties in the analyses. These were discussed in detail in Ref. [64]. The largest uncertainty comes from astrophysics because the trapping rate $\Gamma_t \propto n_x/\bar{v}$. If we naively take the limits $1.75 \leqslant v_{100} \leqslant 2.6$, $0.8 \leqslant r_{10} \leqslant 1$, and $0 \leqslant a_{10} \leqslant 2.5$ ($\bar{v} = \sqrt{3}/2v_r$)

$$0.14 \leqslant \left(\frac{n}{0.3\ \mathrm{cm}^{-3}}\right)\left(\frac{300\ \mathrm{km\ s}^{-1}}{\bar{v}}\right) \leqslant 2.3 \tag{63}$$

TABLE I

Photinos

$m_{\tilde{\gamma}}$ (GeV)	r (Case 1)	r (Case 2)	r (Case 3)
4	0.11	0.014	0.097
6	0.032	0.015	0.077
10	0.0067	0.016	0.054
20	—	0.0066	0.013
40	—	0.0023	0.0023
limit	0.14	0.024	0.040

TABLE II

Higgsinos

$m_{\tilde{H}}$ (GeV)	r (Case 1)	r (Case 2)	r (Case 3)
6	0.045	0.0091	0.052
10	0.0084	0.010	0.040
20	—	0.0057	0.012
40	—	0.0023	0.0025
limit	0.14	0.024	0.040

TABLE III

Majorana neutrinos

m_{ν_M} (GeV)	r (Case 1)	r (Case 2)	r (Case 3)
6	0.055	0.011	0.064
10	0.023	0.028	0.11
20	0.0058	0.051	0.11
40	—	0.078	0.082
limit	0.14	0.024	0.040

implying the signal could be lowered as much as by a factor of 7. However Flores [73] has argued that a conservative lower limit on ϱ_x might be 0.2 GeV cm^{-3} implying a smaller uncertainty. Other factors to be taken into account are fragmentation effects [66], and couplings to protons [67].

TABLE IV

Dirac neutrinos

m_{ν_D} (GeV)	r (Case 1)	r (Case 2)	r (Case 3)
4	—	0.022	0.087
6	—	0.044	0.12
10	—	0.10	0.17
20	—	0.30	0.24
40	—	0.70	0.28
limit	0.14	0.024	0.040

TABLE V

Scalar electron neutrinos

$m_{\tilde{\nu}_e}$ (GeV)	r (Case 1)	r (Case 2)	r (Case 3)
4	—	—	3.7
6	—	—	6.3
10	—	—	11.
20	—	—	20.
40	—	—	26.
limit	0.14	0.024	0.040

TABLE VI

Scalar muon neutrinos

$m_{\tilde{\nu}_\mu}$ (GeV)	r (Case 1)	r (Case 2)	r (Case 3)
4	—	0.50	1.4
6	—	1.3	2.3
10	—	3.8	4.1
20	—	13.	9.5
40	—	35.	18.
limit	0.14	0.024	0.040

Keeping in mind the uncertainties, we can obtain the following limits on cold dark matter candidates with $m_x \gtrsim 3$ GeV:

photinos:	$m_{\tilde{\gamma}} \gtrsim 15$ GeV
Higgsinos:	$m_{\tilde{H}} \gtrsim 10$ GeV
Majorano Neutrinos:	excluded as dark matter
Dirac Neutrinos:	excluded as dark matter
e, μ sneutrinos:	excluded.

One should be very aware however that with the exception of the sneutrinos all limits are within a factor of 2–3 of being acceptable so conservatively we can conclude only the sneutrinos are excluded.

LECTURE III

In this last lecture, I will look at seven instances in which models of particle physics containing a scalar potential which along some direction (linear combination of scalar fields) is flat, $V(\phi) \simeq 0$. I will then examine the cosmological consequence of these potentials.

1. Inflation

As was described in the first lecture, the new inflationary scenario [22] requires a flat potential in order to have a sufficient period of exponential expansion. Though, we clearly do not want $V(\psi) = 0$, inflation requires $\partial V/\partial\psi$ and $\partial^2 V/\partial\psi^2$ to be very small for some extended region in ψ-space. Using the condition (46) and approximating the duration of inflation by

$$\tau \sim \psi/\dot{\psi} \sim H/(\partial^2 V/\partial\psi^2) \tag{64}$$

so that the flatness condition is

$$\partial^2 V/\partial\psi^2 \lesssim H^2/65. \tag{65}$$

As was shown earlier, a potential of the form $V(\psi) = m^4 P(\psi)$ with $m^2/M_P^2 \simeq 10^{-7}$ easily satisfies (65) and yields density fluctuations with $\delta\varrho/\varrho \sim 10^{-4}$ as well.

After the period of exponential expansion, the inflation ψ will begin to roll down to the global minimum at $\langle\psi\rangle = v \simeq M_P$. But because ψ is weakly coupled (usually only gravitationally coupled [29]) ψ oscillates about the minimum until its decay rate becomes comparable to the Hubble rate H. During this period the energy density in the scalar field oscillations is

$$\varrho_\psi = m_\psi^2 \langle\psi\rangle_0^2 (R_\psi/R)^3 \tag{66}$$

where $m_\psi \sim m^2/M_P \sim 10^{-7}\, M_P$, $\langle\psi\rangle_0 \sim M_P$ is the initial amplitude of oscillation and R_ψ is the cosmological scale factor when the oscillations begin. During this period, $H \sim \varrho_\psi^{1/2}/M_P$. For a gravitationally coupled inflaton we might expect a decay rate $\Gamma_\psi \sim m_\psi^3/M_P^2$, so that ψ decays when $R = R_{d\psi}$ given by

$$R_\psi/R_{d\psi} \simeq m_\psi^{4/3}/\langle\psi\rangle_0^{2/3} M_P^{2/3} \tag{67}$$

and the energy density at that time is $\varrho_\psi \sim m_\psi^6/M_P^2$ and the reheat temperature is [26] (assuming rapid thermalization which is in general not true) $T_R \sim \varrho_\psi^{1/4} \sim m_\psi^{3/2}/M_P^{1/2} \sim 10^8$ GeV. This is low enough to make baryogenesis difficult [74] (though not impossible).

2. *Axions*

Axions are pseudo-Goldstone bosons which arise in solving the strong *CP* problem [75, 76] via a global U(1) Peccei-Quinn symmetry. The invisible axion [76] is associated with the flat direction of the spontaneously broken PQ symmetry. Because the PQ symmetry is also explicitly broken (the *CP* violating $\theta F\tilde{F}$ coupling is not PQ invariant) the axion picks up a small mass similar to pion picking up a mass when chiral symmetry is broken. We can expect that $m_a \sim m_\pi f_\pi/f_a$ where f_a, the axion decay constant, is the vacuum expectation value of the PQ current. If we write the axion field as $a = f_a\theta$, near the minimum, the potential produced by QCD instanton effects looks like

$$V(\theta) = \tfrac{1}{2}\, m_a^2\theta^2 f_a^2. \tag{68}$$

Schematically this potential is shown in Fig. 7. In the absense of any *CP* violating effects $\theta_0 = 0$.

The equations of motion for θ can be written as

$$\ddot{\theta} + 3H\dot{\theta} = -m_a^2\theta. \tag{69}$$

For $H \gg m_a$, $\theta \simeq$ constant while for $H < m_a$, θ begins to oscillate about θ_0. The energy density in these oscillations may be the dominant contribution to the mass density of the Universe today [77]. Depending somewhat on the value of f_a, oscillations begin when $T \sim T_i \sim 1$ GeV. The energy density is $\varrho_a = V(\theta)$ where it is understood that m_a is also a temperature dependent quantity. Entropy conservation tells us that $m_a\theta^2$ is constant. So that we can write

$$\varrho_a = \tfrac{1}{2}\, m_a m_a(T_i)\theta_i^2(R_i/R)^3 f_a^2, \tag{70}$$

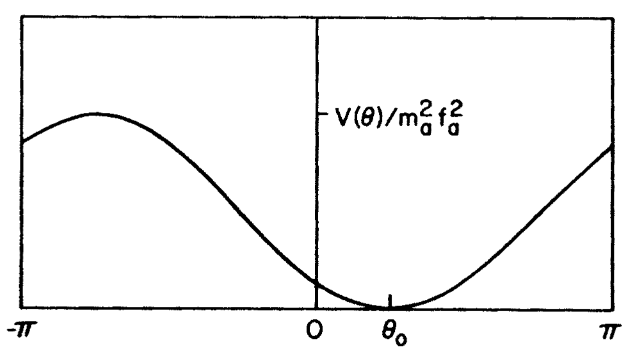

Fig. 7. Schematic drawing of the axion potential in units of $m_a^2 f_a^2$ produced by QCD instaton effects

where $\theta_i \sim O(1)$ is the initial value of θ and R_i is the value of the scale factor when oscillations begin. Taking $m_a(T_i) \sim H(T_i) \sim T_i^2/M_P$, $R_i/R \sim T_0/T_i$ where $T_0 = 2.7$ K, and $m_a \sim 7\times 10^7$ eV (1 GeV/f_a) [78], we find [77]

$$\varrho_a \sim m_a f_a^2 \theta_i^2 T_0^3/M_P T_i \sim 10^{-17}(f_a/1\text{GeV})\ \text{GeV/cm}^3 \tag{71}$$

so that $\Omega \lesssim 1$ implies that

$$f_a \lesssim 10^{12}\ \text{GeV}. \tag{72}$$

It is interesting to note that the recent supernova SN1987a places the strongest constraint from below on f_a [79]

$$f_a \gtrsim 11^{11}\ \text{GeV}. \tag{73}$$

3. *Supergravity breaking*

In supergravity theories [80], unbroken supersymmetry [81] guarantees a degeneracy between fermion-boson (particle-sparticle partners) masses. Supersymmetry breaking lifts this degeneracy so that

$$\Delta m^2 \sim \varepsilon\mu^2, \tag{74}$$

where μ is some supergravity breaking scale and ε is the coupling of the gravitino to the matter and gauge supermultiplets. Typically we expect

$$\varepsilon \sim m_{3/2}/M_P \sim \mu^2/M_P^2 \tag{75}$$

so that

$$\Delta m^2 \sim m_{3/2}^2 \sim \mu^4/M_P^2. \tag{76}$$

Supergravity can be broken by the superhiggs mechanism [82–84]. This can be achieved by introducing a single chiral superfield containing a scalar field z — the Polonyi [83] field and its superpartner the goldstino which gets eaten by the gravitino. The scalar potential in $N = 1$ supergravity can be written in general as [84]

$$V = e^G[G_i(G^{-1})^i_j G^j - 3], \tag{77}$$

where G is the Kahler potential (the metric in field space) such that kinetic energy terms in the Lagrangian are determined by

$$\mathscr{L}_{\text{K.T.}} = -\tfrac{1}{2}\, G^i_j(\partial_\mu \phi^j)(\partial^\mu \phi_i^*) \tag{78}$$

and $G_i = \partial G/\partial\phi^i$, $G^j = \partial G/\partial\phi_j^*$ etc. In minimal $N = 1$ supergravity $G^i_j = \delta^i_j$ and we can write

$$G = \phi^i\phi_i^* + \ln|F|^2, \tag{79}$$

where $F(\phi)$ is the superpotential.

Using Eq. (79) the scalar potential becomes [84]

$$V = e^{\phi_i^* \phi^i}[|F\phi^i + \phi_i^* F|^2 - 3|F|^2], \tag{80}$$

in units where $M_P^2/8\pi = 1$.

The simplest model for breaking supergravity utilizes the superpotential

$$F(z) = \mu^2(z + \Delta), \tag{81}$$

so that

$$V(z, z^*) = \mu^4 e^{|z|^2}[(1 - 3\Delta^2) - 2(z + z^*)\Delta$$
$$+ |z|^2(-1 + \Delta^2 + (z + z^*)\Delta + |z|^2)]. \tag{82}$$

Along the direction $z = z^*$, V has a minimum at $\langle z \rangle = v = (\sqrt{3} - 1)$ and at the minimum $V(v) = 0$ if $\Delta = (2 - \sqrt{3})$. The two scalar fields have masses $m_A^2 = 2\sqrt{3}\, m_{3/2}^2$ and $m_g^2 = 2(2 - \sqrt{3})\, m_{3/2}^2$ where the gravitino mass is

$$m_{3/2}^2 \equiv e^G = \mu^4 e^{(\sqrt{3} - 1)^2}. \tag{83}$$

With $\mu \sim 10^{10}$ GeV, $m_{3/2} \sim 100$ GeV [85]. Clearly this potential is very flat since $V(0)^{1/4} \sim \mu \ll v$. This potential is shown in Fig. 8.

In the early universe the evolution [86] of z is determined by the equation of motion

$$\ddot{z} + 3H\dot{z} = -\partial V/\partial z, \tag{84}$$

when $H > \mu^2/M_P \sim m_{3/2}$, z is constant and when $H < \mu^2/M_P$, z begins to oscillate about the minimum at v. The initial value of $\langle z \rangle$ might be determined by thermal effects or by fluctuations during inflation. In the former, the minimum at finite temperature is at $\langle z \rangle_T \neq v$. In the latter because the mass of the Polonyi field z is so small, fluctuations in $\langle z^2 \rangle$ are induced during inflation [87]. When $m_z^2 \ll H^2$, $\langle z^2 \rangle \sim H^3 t$ up to a limiting value $\langle z^2 \rangle = 3H^4/8\pi^2 m_z^2$ (in our case it would be much smaller since m_z^2 is very large when

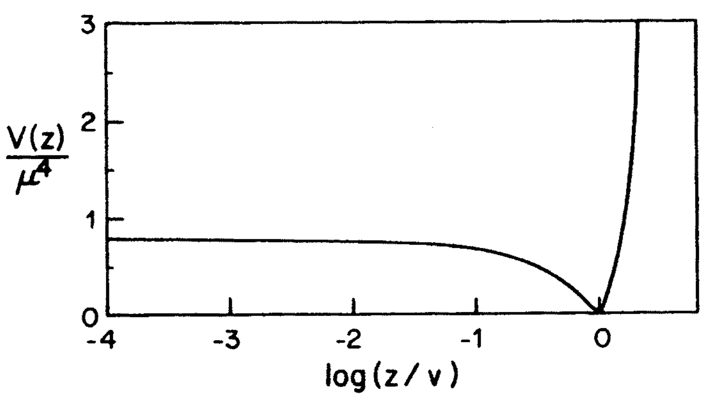

Fig. 8. The scalar potential for the Polonyi field z giving rise to the SuperHiggs mechanism breaking supergravity

$\langle z\rangle \gtrsim v$). In either case we expect that $\langle z\rangle_i \simeq \langle z^2\rangle^{1/2} \sim M_P$ but not at v so that $v-\langle z\rangle_i \sim M_P$ as well.

When z begins oscillating at $T\sim\mu$ the energy density is $\varrho\sim m_z^2(\Delta z)^2\sim\mu^4$ and subsequently $\varrho\sim\mu T^3$. As with the inflaton, oscillations continue until z decays at $H\sim\mu^{1/2}T^{3/2}/M_P\sim\Gamma_z\simeq m_z^3/M_P^2\sim\mu^6/M_P^5$, or at $T_{dz}\sim\mu^{11/3}/M_P^{8/3}$. For $\mu\sim 10^{10}$ GeV, $T_{dz}\sim 10^{-2}$ eV ($T_0 = 2.7$ K $\sim 2\times10^{-4}$ eV) and "reheats" to $T_R\sim\varrho_{osc}^{1/4}(T_d)\sim\mu^3/M_P^2 \sim 1$ keV. The problem with that is that nucleosynthesis would have taken place during the oscillations which is characteristic of a matter dominated expansion rather than a radiation dominated one. Because of the difference in the expansion rate the abundances of the light elements would be greatly altered [88]. Even more problematic is the entropy release due to the decay of these oscillations. The entropy increase is [86] is

$$S_f/S_i \simeq (T_R/T_{dz})^3 \sim (M_P/\mu)^2 \sim 10^{16}. \tag{85}$$

Far too much to understand the present value of η in light of nucleosynthesis and baryosynthesis.

The presence of a flat direction for the breaking of supergravity clearly has an important impact, albeit a negative one, on the early Universe. The simplest solution, to raise the scale of μ, is not acceptable in this context because it would destroy the mass hierarchy which requires $\Delta m^2\sim\mu^4/M_P^2 \lesssim (1\ \mathrm{TeV})^2$. In no-scale supergravity [89] models and in superstring theories [90], we expect the scale for μ to be larger [91] $\mu\sim 10^{-4}\,M_P$ but still leaving us with an entropy increase of $\sim 10^8$. If the baryon asymmetry is produced along the lines discussed in the next section, the Polonyi problem [86] can be solved [92].

4. *Flat directions in the supersymmetric standard model and baryogenesis*

In addition to the flat direction associated with the small scale of supersymmetry breaking, the standard supersymmetric model also contains numerous other flat directions as well. Along these directions squarks and sleptons have in general non-zero vacuum expectation values. Supersymmetry breaking lifts the degeneracy and the "symmetry restoration" can lead to a sizeable baryon asymmetry in the context of a GUT [93].

Let us consider the superpotential for the standard model

$$f = \lambda_1 \bar{H} Q d^c + \lambda_2 H Q u^c + \lambda_e \bar{H} L e^c + \lambda_4 m H \bar{H}, \tag{86}$$

where Q and L represent $SU(2)_L$ doublets of quarks and leptons, u^c, d^c, e^c are the SU(2) singlets and H and $\bar{H}$ are doublet Higgses. These are just the Yukawa mass terms for quarks and charged leptons plus a mixing term for H and $\bar{H}$. The scalar potential is

$$V = |F|^2 + |D|^2 = \sum_i |\partial f/\partial\phi_i|^2 + \frac{g^2}{2}\sum_a (\phi_i^* T_j^{ai}\phi^j)^2, \tag{87}$$

where g is the gauge coupling and T_j^i is a generator of the gauge group. A flat direction can be constructed by taking some linear combination of squark and slepton fields such that $V = 0$. For example, taking [93] $u_3^c = a$, $s_c^2 = a$, $-u_1 = v$, $b_1^c = e^{i\alpha}(v^2+a^2)^{1/2}$ and $\mu^- = v$ with $H = \bar{H} = 0$, is both F and D flat under $SU(3)_C \times SU(2)_L \times U(1)_Y$. Another simple example is [94] $d_1^c = s_2^c = t_3^c = v$. In both cases v and a are arbitrary. Supersymmetry breaking destroys the degeneracy so that scalars typically pick up masses of the order $\tilde{m}^2 \lesssim (1\ \text{TeV})^2$ (or order μ^4/M_P^2 in the notation of the previous section). The potential for this combination of sfermion fields ϕ is shown in Fig. 9. $\langle\phi\rangle_0$ represents some initial value for the expectation value of ϕ.

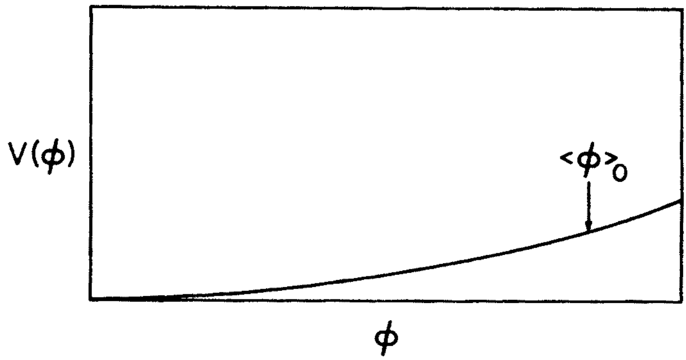

Fig. 9. Schematic view of the scalar potential along a flat direction after supersymmetry breaking effects are included. $\langle\phi\rangle_0$ is the expected initial value of $\langle\phi\rangle$ after inflation

The fact that at early times when the effects of supersymmetry can be neglected, sfermions (charged and colored) have non-zero vacuum expectation values implies that all gauge symmetries were broken, i.e. $M_v \sim g\langle\phi\rangle$ for vector masses. It is interesting that supersymmetry breaking actually restores the gauge symmetries by picking out the ground state at $\langle\phi\rangle = 0$.

When $\langle\phi\rangle \neq 0$, in a GUT such as SU(5), there exists operators V, of the form $qqql$ such that $\langle V\rangle \neq 0$. We can then define a baryon number per particle as [93]

$$B = \text{Im}\,\langle V\rangle/\tilde{m}^2\langle\phi\rangle_0^2 \sim \varepsilon\langle\phi\rangle_0^2/M_G^2, \tag{88}$$

where ε is a measure of CP violation. When the expansion rate of the Universe, H falls below $\tilde{m}$, ϕ begins oscillations about the origin so that the baryon density is

$$n_B \simeq Bn_\phi \simeq B\tilde{m}\langle\phi\rangle^2 \simeq B\tilde{m}\langle\phi\rangle_0^2(R_\phi/R)^3, \tag{89}$$

where $R = R_\phi$ when oscillations begin. As before we can compute the final asymmetry when the oscillations decay when $\Gamma_\phi \sim \tilde{m}^3/\langle\phi\rangle^2 \sim H \sim \tilde{m}\langle\phi\rangle/M_P$. The final baryon to photon ratio is [93]

$$\frac{n_B}{n_\gamma} \sim \varepsilon\,\frac{\langle\phi\rangle_0^2}{M_G^2}\left(\frac{M_P}{\tilde{m}}\right)^{1/6}, \tag{90}$$

which as one can see could be $O(1)$!

Because of the fact that this asymmetry is produced late, $T_R \sim 10^4$ GeV, there are very few dissipative processes to damp this asymmetry before nucleosynthesis [92]. Is this asymmetry embarrassingly large?

Normally we consider the origin of $\langle\phi\rangle_0 \neq 0$ to be due to inflaton, via the scalar field fluctuations [87] discussed previously. That is, $\langle\phi\rangle_0^2 \simeq \langle\phi^2\rangle$ which grows as $H^3 t$ during inflation. We can ask therefore, what is $\langle\phi\rangle_0$ in a typical inflation model and can any of the entropy produced by inflation be used to damp the baryon asymmetry [95]. For our inflaton potential $V(\psi) = m^4 P(\psi)$, recall that the duration of inflation was given by (64) $\tau \sim H/(m^4/M_P^2) \sim M_P^2/m^2$ so that $\langle\phi^2\rangle \sim H^3\tau \sim m^2$ and we should take $\langle\phi\rangle_0^2 \sim \langle\phi^2\rangle \sim 10^{-7} M_P^2$. Furthermore, after inflation, the expansion rate is dominated by inflaton oscillations so that really sfermion decay is determined by $\Gamma_\phi \sim \tilde{m}^3/\langle\phi\rangle^2 \sim H \sim m\langle\psi\rangle/M_P$ with $\langle\psi\rangle_0 \sim M_P$. Then the baryon asymmetry becomes [95]

$$\frac{n_B}{n_\gamma} \sim \frac{\varepsilon\langle\phi\rangle_0^4 m_\psi^{3/2}}{M_G^2 \tilde{m} M_P^{5/2}} . \tag{91}$$

For, $\varepsilon \sim 10^{-3}$, $\langle\phi_0\rangle^2 \sim 10^{-7} M_P^2$, $M_G \sim 10^{-1} M_P$, $m_\psi \sim 10^{-7} M_P$ and $\tilde{m} \sim 10^{-16} M_P$ we find $n_B/n_\gamma \sim 3\times 10^{-10}$ in remarkable agreement with the desired value from nucleosynthesis.

5. Gauge symmetry breaking in superstring models

In this section I consider the breaking of a gauge symmetry in 10-dimensional heterotic superstring theories [90, 96] based on Calabi-Yau compactification [97]. In such theories, the resulting gauge group subsequent to compactification, with SU(3) holonomy [98] is either a rank-5 or rank-6 subgroup of E_6. Recall that the standard model is rank-4. For example the model based on the $CP^3 \times CP^3$ Calabi-Yau manifold yields $SU(3)^3$ as a rank-6 gauge group [99] and one based on CP^7 yields $SU(3)\times SU(2)\times U(1)^2$ as a rank-5 gauge group [100].

The matter fields are contained in the **27** (or $\overline{\mathbf{27}}$) of E_6 which has the following decomposition with respect to SO(10) and SU(5) respectively

$$\mathbf{27} = (\mathbf{16}+\mathbf{10}+\mathbf{1}) = (\mathbf{10}+\bar{\mathbf{5}}+\mathbf{1}+\mathbf{5}+\bar{\mathbf{5}}+\mathbf{1}) \tag{92}$$

and the most general superpotential is of the form

$$F = \lambda_1 HQu^c + \lambda_2 \bar{H}Qd^c + \lambda_3 \bar{H}Le^c + \lambda_4 H\bar{H}N + \lambda_5 NDD^c$$
$$+ \lambda_6 DQQ + \lambda_7 D^c u^c d^c + \lambda_8 D^c QL + \lambda_9 Du^c e^c + \lambda_{10} Dd^c \nu^c + \lambda_{11} HL\nu^c, \tag{93}$$

where H, $\bar{H}$, Q, L, u^c, d^c, e^c have been defined above and D, D^c are new charge $\pm\frac{1}{3}$ color triplets and N and ν^c are standard model singlets (ν^c is like a right handed neutrino). If all of these terms were present, they would result in serious problems [101, 102]. For example, the presence of λ_{11} leads to an unacceptably large Dirac mass for neutrinos and the presence of λ_6 and λ_{10} leads to fast proton decay through $p \to \pi^+\nu^c$. In the above superpotential the first three terms represent the standard model couplings. The λ_4 term supplies a necessary

H $\bar{\mathrm{H}}$ mixing and λ_5 supplies a mass to the D-quark $m_D \simeq \langle N \rangle \simeq 1$ TeV. By inclusion of a simple Z_2 discrete symmetry $(D, D^c, \nu^c) \to -(D, D^c, \nu^c)$ the unwanted couplings λ_6, λ_7, λ_8, λ_9 and λ_{11} can be made to vanish. This is characteristic of all inspired models where discrete symmetries must be imposed to delete unwanted Yukawa couplings.

In a rank-5 model [101], only one of the two standard model singlets ν^c or N need pick up a vacuum expectation value. As noted above a vev for N produces H$\bar{\mathrm{H}}$ mixing and D masses. This phase transition has been shown [92] to be cosmologically safe from the point of view of entropy production. In a rank-6 model, both ν^c *and* N need to pick up vacuum expectation values. It has been shown [103] that without an intermediate scale (e.g. $\langle \nu^c \rangle \simeq O(10^{10}$ GeV) the addition of a ν^c vev produces problems for the down quark mass matrix, flavor changing neutral currents and lepton number violation.

There are however severe problems which can arise in models with intermediate scales. Intermediate scales arise along flat directions of a scalar potential of the form

$$V(\phi, T) = (-\tilde{m}^2 + T^2)\phi^2 + \phi^{4+n}/M_c^n, \tag{94}$$

where ϕ might be ν^c for example, $m_\phi^2 \simeq \tilde{m}^2$ and $M_c \sim 10^{-1}\, M_P$ is the compactification scale. The zero temperature potential has a minimum at

$$\langle \phi \rangle = M_I \simeq (\tilde{m}^2 M_c^2)^{1/(2+n)}, \tag{95}$$

so that for $n = 2$, $M_I \simeq 10^{10}$ GeV. Larger intermediate scales are possible along flatter directions (i.e. with $n > 2$). See Fig. 10.

Problems for intermediate scale models can arise in many ways. First, in many models, baryon decay can be mediated by dimension-five operators due to the exchange of particles weighing $O(M_I)$. The stability of the proton suggests that $M_I \gtrsim 10^{16}$ GeV [105]. Second, in models with several mirror generations, unless $M_I \gtrsim 10^{16-17}$ GeV the evolution of gauge couplings is rapid enough to preclude a perturbative regime at lowe energies. Third, unless $M_I \lesssim 10^{15}$ GeV or there exist non-zero scalar masses at the compactification scale, it is very unlikely that the renormalization group equation evolution of scalar masses would be strong enough to drive $m_\phi^2 < 0$ to generate the intermediate scale. Finally, unless

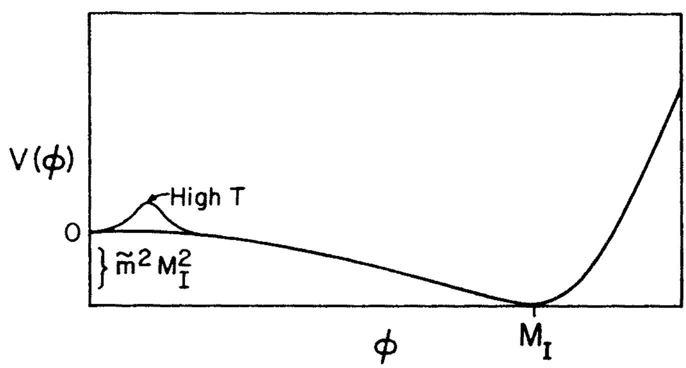

Fig. 10. Schematic view of the scalar potential along a flat direction breaking an extended gauge symmetry down to the standard model, in superstring theories with an intermediate scale M_I

$M_I < 10^7$ GeV (10^{10-12} GeV depending on assumptions about inflation and baryosynthesis) the entropy produced in this phase transition is too large [106] thereby diluting the baryon asymmetry of the Universe. I will discuss only this last problem in more detail.

Using the scalar potential in equation (94) we see that the phase transition $\langle\phi\rangle = 0$ to $\langle\phi\rangle = M_I$ in the early Universe takes place at $T \simeq \tilde{m} \leqslant 1$ TeV. For $T \leqslant \tilde{m}$, the ϕ field, as in the many cases before, oscillates about the minimum at M_I with an energy density $\varrho_\phi \sim \tilde{m}^2 M_I^2 (R_\phi/R)^3$ and decays when $\Gamma_\phi \sim \tilde{m}^3/M_I^2 \sim H$. The entropy produced after ϕ decay is [106]

$$\Delta = S_f/S_i \simeq M_I^3/m_\phi^{5/2} M_c^{1/2}. \tag{96}$$

Unless $M_I \leqslant 10^7$ GeV, $S_f/S_I \geqslant 10^6$.

This bound can be relaxed somewhat by considering the effects of inflation. Symmetry breaking during inflation effectively induces ϕ oscillations earlier ($T > \tilde{m}$) so that if the reheating temperature after inflation does not restore the symmetry and $T_R < (\tilde{m}M_I)^{1/2}$ then $S_f/S_i < 10^6$ if $M_I < 10^{10}$ GeV [104]. This bound can be further relaxed if the baryon asymmetry were produced after the phase transition such as in the mechanism described above [93]. But this only raises the limit of M_I to $M_I \leqslant 10^{12}$ GeV.

A promising alternative to the 10-dimensional models is the 4-dimensional formulation of superstring theory [107]. One example which has been examined [108] in detail is the model based on the gauge group $SU(5) \times U(1)$. The model is attractive from a model--building point of view in that it does not require an adjoint to break SU(5), it is done with a **10**-plet. There is also a very natural doublet-triplet separation mechanism and a see-saw for neutrino masses. Although $\Delta B \neq 0$ interactions require a scale $M_I \simeq 10^{16}$ GeV, strong coupling effects become important at $T \simeq \Lambda_5 \simeq 10^{10}$ GeV and can avoid the entropy problem normally associated with intermediate scale models [109].

6. *Inflation and the generation of an electric charge asymmetry*

In this section I will briefly describe a problem [110] that arises, during inflation in the presence of the same flat direction used to generate the baryon asymmetry in Section 4. As was noted above during and just after inflation, all gauge symmetries are broken by the expectation value for the sfermions ϕ. In particular, electric charge is violated and the photon has a mass $\sim e\langle\phi\rangle_0$. Because of this, an electric charge asymmetry may have been generated after inflation.

Because our gauge group contains no operators which explicitly violate electric charge, we can not generate a charge asymmetry in the same way as we argued for baryons. But other processes, in particular inflaton decay, now has charge violating modes which are in all likelihood CP violating as well. Consider a piece of the lagrangian which gravitationally couples the inflaton to matter

$$\mathscr{L} \ni \frac{m_\psi}{M_P} h\psi\bar{H}\varphi d^c, \tag{97}$$

where h is some Yukawa coupling.

If for example we consider a flat direction containing d^c then by shifting the fields, we get a charge violating interaction of the form $(m_\psi/M_P)\langle\phi\rangle h\psi\bar{H}Q$. The branching ratio for the charge violating mode is

$$\Gamma_{\Delta Q}/\Gamma_\psi = h^2\langle\phi\rangle^2/m_\psi^2, \tag{98}$$

where $\langle\phi\rangle^2 = \langle\phi\rangle_0^2(R_\phi/R)^3$ is to be evaluated at the time the inflaton ψ decays. The net charge asymmetry is now [110]

$$\frac{n_Q}{n_\gamma} \simeq \varepsilon' h^2 \frac{\langle\phi\rangle_0^2}{\tilde{m}^2}\left(\frac{m_\psi}{M_P}\right)^{9/2}, \tag{99}$$

where ε' is a measure of the CP violation in the inflaton decay. When compared to the baryon asymmetry, it is found that $n_Q/n_B \sim 10^{-4}$. Limits due to electrostatic repulsions require $n_Q/n_B \lesssim 10^{-18}$ so we are off by some 14 orders of magnitude! Limits from the propagation of cosmic rays imply further that [111] $n_Q/n_B \lesssim 10^{-29}$.

There are not many attractive resolutions to this problem: 1) An unknown suppression of by 14 orders of magnitude in the CP violating phase; 2) Reduce $\langle\phi\rangle_0$ (which also reduces n_B/n_γ), but it is not clearly why $\langle\phi\rangle_0$ would be smaller than its value obtained by fluctuations during inflation; 3) Reduce m_ψ (this would also reduce n_B/n_γ) but this would lower the magnitude of $\delta\varrho/\varrho \simeq 10^3\, m_\psi/M_P$, and thus an additional source of density fluctuations would be needed to produce galaxies; 4) Eliminate the unwanted couplings, but they are already gravitational. The best solution seems to be to consider residual non-renormalization terms in the potential due to integrating out massive fields in a GUT [94], which would reduce $\langle\phi\rangle_0$ and leave $n_Q/n_B \simeq 10^{-23}$ — somewhere between the two observational limits. This clearly imposes a severe problem for inflationary theories in models with flat directions such as the standard supersymmetric model.

7. *Initial conditions for chaotic inflation*

Finally the existence of these same flat directions imposed strong constraints [112] on the initial conditions for chaotic inflation [38]. Models of chaotic inflation are based on very simple potentials such as $V(\psi) = \frac{1}{2} m_\psi^2\psi^2$ or $\frac{1}{4}\lambda\psi^4$ (here I will only consider the former) and the initial conditions for $\langle\psi\rangle$ are fixed by setting $V(\psi) \simeq M_P^4$. The magnitude of density perturbations is just $\delta\varrho/\varrho \simeq 10\, m_\psi/M_P$ so that we take $m_\psi \sim 10^{-5}\, M_P$ in this case and $\langle\psi\rangle_0 \simeq 10^5\, M_P$. The duration of inflation is more than sufficient $H\tau \sim \langle\psi\rangle_0^2/M_P^2 \sim 10^{10}$.

The chaotic condition $V \simeq M_P^4$, however should also be applied to other scalar fields as well. In particular, along the sfermion flat direction this would imply that $\langle\phi\rangle_0 \simeq 10^{16}\, M_P$. The problem [112] that arises is that density fluctuations are fixed by the last field to inflate which in this case would be ϕ not ψ and hence $\delta\varrho/\varrho \simeq 10^{-15}$, far too small to have any relevance for galaxy formation. Indeed even if $\langle\phi\rangle_0 = 0$ initially, scalar field fluctuations [87] drive $\langle\phi^2\rangle$ to $H^3\tau$. When $\langle\psi\rangle_0 \gtrsim 10^3\, M_P$, $\langle\phi^2\rangle$ fluctuations are large enough so that ϕ is again the last to inflate. In Fig. 11, the allowable parameter

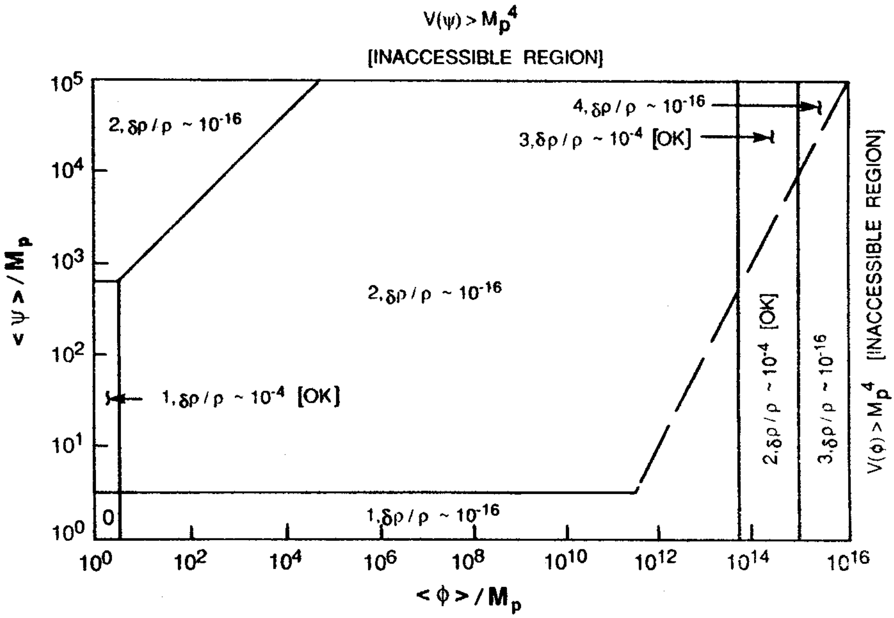

Fig. 11. The magnitude of energy density fluctuations produced and the number of inflating epochs as a function of the sfermion $\langle\phi\rangle$ and inflation $\langle\psi\rangle$ initial values, in a chaotic inflationary scenario

space for $\langle\phi\rangle_0$ and $\langle\psi\rangle_0$ is shown in order to have $\delta\varrho/\varrho \simeq 10^{-4}$. Either $\langle\phi\rangle_0 \lesssim 3\,M_P$ and $3M_P \lesssim \langle\psi\rangle_0 \lesssim 730\,M_P$ or $\langle\psi\rangle_0 \lesssim 10^5\,M_P$ and $6\times10^{13}\,M_P \lesssim \langle\phi\rangle_0 \lesssim 8\times10^{14}\,M_P$. Although a clear non-negligible area for $\delta\varrho/\varrho \simeq 10^{-4}$ is present, complete chaotic conditions may very well result in very little structure.

To summarize, I have tried to show that the existence of flat directions in models of particle physics which appear under almost any circumstance can have a profound impact on cosmology and the evolution of the early Universe.

This work was supported in part by DOE grant DE-AC02-83ER-40105 and by a Presidential Young Investigator Award.

REFERENCES

[1] For a more thorough treatment see: S. Weinberg, *Gravitation and Cosmology*, J. Wiley and Sons, New York 1972.

[2] See e.g., G. A. Tammann, A. Sandage, A. Yahil, *Physical Cosmology*, R. Balian, J. Audouze, D. N. Schramm (eds.), North Holland Pub. Co., Amsterdam 1980.

[3] A. A. Penzias, R. W. Wilson, *Ap. J.* **142**, 419 (1965).

[4] G. Gamow, *Phys. Rev.* **70**, 572 (1946); R. A. Alpher, H. Bethe, G. Gamow, *Phys. Rev.* **73**, 803 (1948); P. J. E. Peebles, *Ap. J.* **146**, 542 (1966); R. V. Wagoner, W. A. Fowler, F. Hoyle, *Ap. J.* **148**, 3 (1967).

[5] D. N. Schramm, R. V. Wagoner, *Ann. Rev. Nucl. Part. Sci.* **27**, 37 (1977); K. A. Olive, D. N. Schramm, G. Steigman, M. S. Turner, J. Yang, *Ap. J.* **246**, 547 (1981); A. Bosegard, G. Steigman, *Ann. Rev. Astron. Astrophys.* **23**, 319 (1985).

[6] J. Yang, M. S. Turner, G. Steigman, D. N. Schramm, K. A. Olive, *Ap. J.* **281**, 493 (1984).
[7] G. Steigman, D. N. Schramm, J. E. Gunn, *Phys. Lett.* **86B**, 202 (1977); J. Yang, D. N. Schramm, G. Steigman, R. T. Rood, *Ap. J.* **227**, 697 (1979).
[8] G. Steigman, K. A. Olive, D. N. Schramm, M. S. Turner, *Phys. Lett.* **176B**, 33 (1986).
[9] E. Witten, *Phys. Rev.* **D30**, 272 (1984).
[10] J. Kapusta, K. A. Olive, *Phys. Lett.* **209B**, 295 (1988).
[11] J. H. Applegate, C. J. Hogan, *Phys. Rev.* **D31**, 3037 (1985); J. H. Applegate, C. J. Hogan, R. J. Scherrer, *Phys. Rev.* **D35**, 1151 (1987); C. Alcock, G. M. Fuller, G. J. Mathews, *Ap. J.* **320**, 439 (1987); G. M. Fuller, G. J. Mathews, C. R. Alcock, *Phys. Rev.* **D37**, 1380 (1988).
[12] H. Reeves, in Proc. Int. Sch. Phys. Enrico Fermi 1987 (in press).
[13] K. A. Olive, *Nature* **330**, 770 (1987).
[14] H. Kurki-Suonio, R. A. Matzner, J. M. Centrella, T. Rathman, J. R. Wilson, *Phys. Rev.* **D38**, 1091 (1988).
[15] A. D. Sakharov, *Zh. Eksp. Teor. Fiz. Pisma. Red.* **5**, 32 (1967).
[16] G. Steigman, *Ann. Rev. Astron. Astrophys.* **14**, 339 (1976).
[17] S. Weinberg, *Phys. Rev. Lett.* **42**, 850 (1979); D. Toussaint, S. B. Treiman, F. Wilczek, A. Zee, *Phys. Rev.* **D19**, 1036 (1979).
[18] J. N. Fry, K. A. Olive, M. S. Turner, *Phys. Rev.* **D22**, 2953 and 2977 (1980).
[19] A. H. Guth, *Phys. Rev.* **D23**, 347 (1981).
[20] S. Coleman, *Phys. Rev.* **D15**, 2929 (1977); C. Callan, S. Coleman, *Phys. Rev.* **D16**, 1762 (1977).
[21] A. H. Guth, E. Weinberg, *Phys. Rev.* **D23**, 826 (1981); and *Nucl. Phys.* **B212**, 321 (1983).
[22] A. D. Linde, *Phys. Lett.* **108B**, 389 (1982); A. Albrecht, P. Steinhardt, *Phys. Rev. Lett.* **48**, 1220 (1982).
[23] W. H. Press, *Phys. Scr.* **21**, 702 (1980); S. W. Hawking, *Phys. Lett.* **115B**, 295 (1982); A. H. Guth, S. Y. Pi, *Phys. Rev. Lett.* **49**, 1110 (1982); A. A. Starobinski, *Phys. Lett.* **117B**, 175 (1982); J. M. Bardeen, P. J. Steinhardt, M. S. Turner, *Phys. Rev.* **D28**, 679 (1983).
[24] S. Coleman, E. Weinberg, *Phys. Rev.* **D7**, 1888 (1973).
[25] J. Ellis, D. V. Nanopoulos, K. A. Olive, K. Tamvakis, *Nucl. Phys.* **B221**, 224 (1983).
[26] D. V. Nanopoulos, K. A. Olive, M. Srednicki, *Phys. Lett.* **127B**, 30 (1983).
[27] J. Ellis, D. Nanopoulos, K. Olive, K. Tamvakis, *Phys. Lett.* **118B**, 335 (1982).
[28] J. Ellis, D. V. Nanopoulos, K. A. Olive, K. Tamvakis, *Phys. Lett.* **120B**, 334 (1983).
[29] D. V. Nanopoulos, K. A. Olive, M. Srednicki, K. Tamvakis, *Phys. Lett.* **123B**, 41 (1983).
[30] R. Holman, P. Ramond, G. G. Ross, *Phys. Lett.* **137B**, 343 (1984).
[31] G. Gelmini, D. V. Nanopoulos, K. A. Olive, *Phys. Lett.* **131B**, 53 (1983).
[32] B. Ovrut, P. Steinhardt, *Phys. Lett.* **133B**, 161 (1983).
[33] G. Mazenko, W. Unruh, R. Wald, *Phys. Rev.* **D31**, 273 (1985).
[34] L. Jensen, K. A. Olive, *Nucl. Phys.* **B263**, 731 (1986).
[35] J. Ellis, K. Enquist, D. V. Nanopoulos, K. A. Olive, M. Srednicki, *Phys. Lett.* **152B**, 175 (1985).
[36] G. Gelmini, C. Kounnas, D. V. Nanopoulos, *Nucl. Phys.* **B250**, 177 (1985).
[37] A. Albrecht, R. Bradenberg, *Phys. Rev.* **D31**, 1225 (1985); G. D. Coughlan, G. G. Ross, *Phys. Lett.* **157B**, 151 (1985); L. Jensen, K. A. Olive, *Phys. Lett.* **159B**, 99 (1985).
[38] A. D. Linde, *Phys. Lett.* **129B**, 177 (1983).
[39] F. Graziani, K. Olynyk, Fermilab preprint 85-175 (1985).
[40] J. Ellis, K. Enqvist, D. V. Nanopoulos, M. Quiros, *Nucl. Phys.* **B277**, 231 (1986); P. Binetruy, M. K. Gaillard, *Phys. Rev.* **D34**, 3069 (1986); F. Graziani, K. A. Olive, *Phys. Rev.* (1988) (in press); C. Kounnas, S. Kalara, K. A. Olive, University of Minnesota preprint UMN-TH-658/88 (1988); S. Kalara, K. A. Olive, University of Minnesota preprint UMN-TH-705/88 (1988).
[41] S. M. Faber, J. J. Gallagher, *Ann. Rev. Astron. Astrophys.* **17**, 135 (1979).
[42] J. Primack, SLAC preprint 3387 (1984).
[43] D. Hegyi, K. A. Olive, *Phys. Lett.* **126B**, 28 (1983); *Ap. J.* **303**, 56 (1986).

[44] J. R. Bond, A. Szalay, *Ap. J.* **274**, 443 (1983).
[45] R. Cowsik, J. McClelland, *Phys. Rev. Lett.* **29**, 669 (1972); A. S. Szalay, G. Marx, *Astron. Astrophys.* **49**, 437 (1976).
[46] G. Steigman, K. A. Olive, D. N. Schramm, *Phys. Rev. Lett.* **43**, 239 (1979); K. A. Olive, D. N. Schramm, G. Steigman, *Nucl. Phys.* **B180**, 497 (1981).
[47] K. A. Olive, M. S. Turner, *Phys. Rev.* **D25**, 213 (1982).
[48] P. Hut, *Phys. Lett.* **69B**, 85 (1977); B. W. Lee, S. Weinberg, *Phys. Rev. Lett.* **39**, 165 (1977).
[49] E. W. Kolb, K. A. Olive, *Phys. Rev.* **D33**, 1202 (1986); E: **34**, 2531 (1986).
[50] R. Watkins, M. Srednicki, K. A. Olive, *Nucl. Phys.* **B**, 1988 (in press).
[51] L. M. Krauss, *Phys. Lett.* **128B**, 37 (1983).
[52] H. Goldberg, *Phys. Rev. Lett.* **50**, 1419 (1983).
[53] L. M. Krauss, *Nucl. Phys.* **B227**, 556 (1983).
[54] J. Ellis, J. Hagelin, D. V. Nanopoulos, K. A. Olive, M. Srednicki, *Nucl. Phys.* **B238**, 453 (1984).
[55] L. E. Ibanez, *Phys. Lett.* **137B**, 160 (1984); J. Hagelin, G. L. Kane, S. Raby, *Nucl. Phys.* **B241**, 638 (1984).
[56] W. H. Press, D. N. Spergel, *Ap. J.* **296**, 67 (1985); A. Gould, *Ap. J.* **321**, 571 (1987).
[57] M. Srednicki, K. A. Olive, J. Silk, *Nucl. Phys.* **B279**, 804 (1987).
[58] J. Silk, K. A. Olive, M. Srednicki, *Phys. Rev. Lett.* **55**, 257 (1985).
[59] K. Freese, D. N. Spergel, W. H. Press, *Ap. J.* **299**, 1001 (1986).
[60] G. Steigman, C. Sarazin, H. Quintana, J. Faulkner, *Ap. J.* **83**, 1050 (1978); D. N. Spergel, W. H. Press, *Ap. J.* **294**, 663 (1985).
[61] K. Greist, D. Seckel, *Nucl. Phys.* **B283**, 681 (1987).
[62] J. Hagelin, K. W. Ng, K. A. Olive, *Phys. Lett.* **180B**, 375 (1986).
[63] T. K. Gaisser, G. Steigman, S. Tilav, *Phys. Rev.* **D34**, 2206 (1986).
[64] K. W. Ng, K. A. Olive, M. Srednicki, *Phys. Lett.* **188B**, 138 (1987).
[65] K. A. Olive, M. Srednicki, *Phys. Lett.* **B205**, 553 (1988).
[66] S. Ritz, D. Seckel, *Nucl. Phys.* **B304**, 877 (1988).
[67] J. Ellis, R. A. Flores, S. Ritz, *Phys. Lett.* **198B**, 393 (1984).
[68] see e.g. D. H. Perkins, *Ann. Rev. Nucl. Part. Sci*, **34**, 1 (1984).
[69] B. Campbell, J. Ellis, K. Enqvist, D. V. Nanopoulos, J. Hagelin, K. A. Olive, *Phys. Lett.* **173B**, 270 (1986).
[70] IMB Collaboration: J. LoSecco et al., *Phys. Lett.* **188B**, 388 (1987).
[71] Kamioka Collaboration: Y. Totsuka, University of Tokyo preprint UT-ICEPP-87-02 (1987).
[72] Frejus Collaboration: B. Kuznik, Orsay preprint LAL 87-21 (1987).
[73] R. Flores, CERN preprint TH-4736 (1987).
[74] J. Ellis, K. Enqvist, G. Gelmini, C. Kounnas, A. Masiero, D. V. Nanopoulos, A. Yv. Smirnov, *Phys. Lett.* **147B**, 27 (1984).
[75] R. D. Peccei, H. R. Quinn, *Phys. Rev. Lett.* **37**, 1440 (1977); *Phys. Rev.* **D16**, 1791 (1977); S. Weinberg, *Phys. Rev. Lett.* **40**, 223 (1978); F. Wilczek, *Phys. Rev. Lett.* **40**, 279 (1978).
[76] J. E. Kim, *Phys. Rev. Lett.* **43**, 103 (1979); M. A. Shifman, A. I. Vainshtein, V. I. Zakharov, *Nucl. Phys.* **B166**, 493 (1980); M. Dine, W. Fischler, M. Srednicki, *Phys. Lett.* **104B**, 199 (1981).
[77] J. Preskill, M. B. Wise, F. Wilczek, *Phys. Lett.* **120B**, 127 (1983); L. F. Abbott, P. Sikivie, *Phys. Lett.* **120B**, 133 (1983); M. Dine, W. Fischler, *Phys. Lett.* **120B**, 137 (1983); M. Turner, *Phys. Rev.* **D33**, 889 (1986).
[78] M. Srednicki, *Nucl. Phys.* **B260**, 689 (1985).
[79] R. Mayle, J. R. Wilson, J. Ellis, K. A. Olive, D. N. Schramm, G. Steigman, *Phys. Lett.* **203B**, 188 (1988) and University of Minnesota preprint UMN-TH-702/88; G. G. Raffelt, D. Seckel, *Phys. Rev. Lett.* **60**, 1793 (1988); M. S. Turner, *Phys. Rev. Lett.* **60**, 1797 (1988); R. P. Brinkman, M. S. Turner; Fermilab preprint (1988).

[80] D. Z. Freedman, P. Van Nieuwenhuizen, S. Ferrara, *Phys. Rev.* **D13**, 3214 (1976); S. Deser, B. Zumino, *Phys. Lett.* **62B**, 335 (1976); P. Van Nieuwenhuizen, *Phys. Rep.* **68C**, 189 (1981).

[81] Y. A. Golfand, E. P. Likhtman, *Pisma Zh. Eksp. Teor. Fiz.* **13**, 323 (1971); D. Volkov, V. P. Akulov, *Phys. Lett.* **46B**, 109 (1973); J. Wess, B. Zumino, *Nucl. Phys.* **B70**, 39 (1974); P. Fayet, S. Ferrara, *Phys. Rep.* **32C**, 249 (1977).

[82] D. V. Volkov, V. A. Soroka, *JETP Lett.* **18**, 312 (1973); S. Deser, B. Zumino, *Phys. Rev. Lett.* **38**, 1433 (1977).

[83] J. Polonyi, Budapest preprint KFKI-1977-93 (1977).

[84] E. Cremmer, B. Julia, J. Scherk, S. Ferrara, L. Giraradello, P. Van Nieuwenhuizen, *Phys. Lett.* **79B**, 23 (1978) and *Nucl. Phys.* **B147**, 105 (1979); E. Cremmer, S. Ferrara, L. Girardello, A. Van Proeyen, *Phys. Lett.* **116B**, 231 (1982) and *Nucl. Phys.* **B212**, 413 (1983).

[85] R. Barbieri, S. Ferrara, C. A. Savoy, *Phys. Lett.* **119B**, 343 (1982).

[86] G. D. Coughlan, W. Fischler, E. W. Kolb, S. Raby, G. G. Ross, *Phys. Lett.* **131B**, 59 (1983).

[87] T. S. Bunsch, P. C. W. Davies, *Proc. R. Soc. Lond.* **A360**, 117 (1978); A. Vilenkin, L. H. Ford, *Phys. Rev.* **D26**, 1231 (1982); A. D. Linde, *Phys. Lett.* **116B**, 335 (1982); A. Vilenkin, *Nucl. Phys.* **B226**, 527 (1983); K. Enqvist, K. W. Ng, K. A. Olive, *Nucl. Phys.* **B303**, 713 (1988).

[88] see e.g. E. W. Kolb, R. Scherrer, *Phys. Rev.* **D25**, 1481 (1982).

[89] E. Cremmer, S. Ferrara, C. Kounnas, D. V. Nanopoulos, *Phys. Lett.* **133B**, 61 (1983); J. Ellis, A. B. Lahanas, D. N. Nanopoulos, K. Tamvakis, *Phys. Lett.* **134B**, 429 (1984); J. Ellis, C. Kounnas, D. V. Nanopoulos, *Nucl. Phys.* **B241**, 406 (1984) and **B247**, 373 (1984); A. B. Lahanas, D. N. Nanopoulos, *Phys. Rep.* **145**, 1 (1987).

[90] M. B. Green, J. H. Schwarz, E. Witten, *Superstring Theory*, Cambridge University Press, Cambridge 1987.

[91] P. Binetruy, M. K. Gaillard, *Phys. Lett.* **168B**, 347 (1986); M. Quiros, *Phys. Lett.* **173B**, 265 (1986); Y. J. Ahn, J. D. Breit, *Nucl. Phys.* **B273**, 75 (1986); P. Binetruy, S. Dawson, I. Hinchliffe, M. K. Gaillard, *Phys. Lett.* **192B**, 377 (1987); J. Ellis, D. V. Nanopoulos, M. Quiros, F. Zwirner, *Phys. Lett.* **180B**, 83 (1986).

[92] J. Ellis, D. V. Nanopoulos, K. A. Olive, *Phys. Lett.* **184B**, 37 (1987).

[93] I. Affleck, M. Dine, *Nucl. Phys.* **B249**, 361 (1985); A. D. Linde, *Phys. Lett.* **160B**, 243 (1985).

[94] K. W. Ng, University of Minnesota preprint UMN-TH-704/88 (1988).

[95] J. Ellis, K. Enqvist, D. V. Nanopoulos, K. A. Olive, *Phys. Lett.* **191B**, 343 (1987).

[96] D. J. Gross, J. Harvey, E. Martinec, R. Rohm, *Phys. Rev. Lett.* **54**, 502 (1985); *Nucl. Phys.* **B256**, 253 (1985) and **B267**, 75 (1986).

[97] P. Candelas, G. T. Horowitz, A. Strominger, E. Witten, *Nucl. Phys.* **B258**, 46 (1985).

[98] E. Witten, *Nucl. Phys.* **B258**, 75 (1985).

[99] B. Greene, K. H. Kirklin, P. J. Miron, G. C. Ross, *Phys. Lett.* **180B**, 69 (1986); *Nucl. Phys.* **B274**, 574 (1986) and **B292**, 606 (1987); see also J. Ellis, K. Enqvist, D. V. Nanopoulos, K. A. Olive, *Nucl. Phys.* **B297**, 103 (1988); S. Kalara, R. N. Mohapatra, *Phys. Rev.* **D36**, 3474 (1987).

[100] J. Ellis, K. Enqvist, S. Kalara, D. V. Nanopoulos, K. A. Olive, *Nucl. Phys.* **B306**, 445 (1988).

[101] J. Ellis, K. Enqvist, D. V. Nanopoulos, F. Zwirner, *Nucl. Phys.* **B276**, 14 (1986) and *Mod. Phys. Lett.* **A1**, 57 (1986).

[102] B. Campbell, J. Ellis, K. Enqvist, M. K. Gaillard, D. V. Nanopoulos, *Int. J. Mod. Phys.* **2A**, 831 (1987).

[103] B. Campbell, J. Ellis, M. K. Gaillard, D. V. Nanopoulos, K. A. Olive, *Phys. Lett.* **180B**, 77 (1987); B. Campbell, K. A. Olive, D. B. Reiss, *Nucl. Phys.* **B296**, 129 (1988).

[104] J. Ellis, K. Enqvist, D. V. Nanopoulos, K. A. Olive, *Phys. Lett.* **188B**, 415 (1987).

[105] B. Campbell, J. Ellis, D. V. Nanopoulos, *Phys. Lett.* **141B**, 229 (1984).

[106] K. Yamamoto, *Phys. Lett.* **168B**, 341 (1986); K. Enqvist, D. V. Nanopoulos, M. Quiros, *Phys. Lett.* **169B**, 343 (1986).

[107] K. S. Narain, *Phys. Lett.* **169B**, 41 (1986); H. Kawai, D. C. Lewellen, H. H. Tye, *Phys. Rev.*

Lett. **57**, 1832 (1986); *Nucl. Phys.* **B288**, 1 (1987); I. Antoniadis, C. Bachas, C. Kounnas, *Nucl. Phys.* **B289**, 87 (1987).
[108] I. Antoniadis, J. Ellis, J. S. Hagelin, D. V. Nanopoulos, *Phys. Lett.* **194B**, 231 (1987); and **205B**, 459 (1988).
[109] B. A. Campbell, J. Ellis, J. S. Hagelin, D. V. Nanopoulos, K. A. Olive, *Phys. Lett.* **197B**, 335 (1987); **200B**, 483 (1988) and **207B**, 451 (1988).
[110] K. Enqvist, K. W. Ng, K. A. Olive, University of Minnesota preprint UMN-TH-656/88 (1988).
[111] S. Orito, M. Yoshimura, *Phys. Rev. Lett.* **54**, 2457 (1985).
[112] F. Graziani, K. A. Olive, University of Minnesota preprint UMN-TH-660/88 (1988).