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PARTICLE PHYSICS AND COSMOLOGY: THE INNER SPACE/OUTER SPACE CONNECTION*

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ABSTRACT

Advances made in elementary particle physics in recent years are helping cosmologists to understand the very early history of the Universe ($t \ll 10^{-2}$ sec). At the same time particle physicists now routinely use the early Universe and various contemporary astrophysical environments (red giants, neutron stars, etc.) as non-traditional HEP laboratories in which physics at very high energies ($\gg 10^3$ GeV) can be probed. In this talk I will review both of these aspects of the 'Inner Space/Outer Space Connection.'

INTRODUCTION

In the past five years or so progress in both elementary particle physics and in cosmology has become increasingly dependent upon the interplay between the two disciplines. On the particle physics side, the $SU(3)_C \times SU(2)_L \times U(1)_Y$ model seems to very accurately describe the interactions of quarks and leptons at energies below, say, 10^3 GeV. At the very least, the so-called standard model is a satisfactory, effective low-energy theory. The frontiers of particle physics now involve energies of much greater than 10^3 GeV--energies which are not now available in terrestrial accelerators, nor are ever likely to be available in terrestrial accelerators. For this reason particle physicists have turned both to the early Universe with its essentially unlimited energy budget (up to 10^{19} GeV) and high particle fluxes (up to 10^{107} cm^{-2} s^{-1}), and to various unique, contemporary astrophysical environments (centers of main sequence stars where temperatures reach 10^8 K, neutron stars where densities reach 10^{14}-10^{15} g cm^{-3}, our galaxy whose magnetic field can impart 10^{11} GeV to a Dirac magnetic charge, etc.) as non-traditional laboratories for studying physics at very high energies and very short distances.

On the cosmological side, the hot big bang model, the so called standard model of cosmology, seems to provide an accurate accounting of the history of the Universe from about 10^{-2} sec after 'the bang' when the temperature was about 10 MeV, until today, some 10-20 billion years after 'the bang' and temperature of about 3 K ($\approx 3 \times 10^{-13}$ GeV). Extending our understanding further back, to earlier times and higher temperatures, requires knowledge about the fundamental particles (presumably quarks and leptons) and their interactions at very high energies. For this reason, progress in cosmology has become linked to progress in elementary particle physics.

In this review I will focus on both aspects of the interplay between particle physics and cosmology. I will begin by briefly reviewing the standard cosmology. I will then discuss the implications of the two most spectacular predictions of grand unified theories (GUTs) (baryon number nonconservation and superheavy magnetic monopoles) for the standard cosmology. As successful as it is, the hot big bang model fails to account for a number of very fundamental cosmological observational facts, including the large-scale isotropy and homogeneity (of the observable Universe), the small-scale inhomogeneity, the near critical expansion rate (or flatness of the Universe), and in the context of GUTs, the fact that we do not live in a 30,000 yr old Universe where the monopole abundance is $O(10^{-3})$ per baryon (corresponding to a flux of $\approx 10^{-4}$ cm^{-2} sr^{-1} s^{-1}!). The inflationary Universe models originally conceived by Guth[1] and developed by Linde[2] and Albrecht and Steinhardt[3] go a long way toward resolving these cosmological conundrums (and may in fact solve them). After discussing inflation, I will address the question of the nature of the dark matter in the Universe. Rotation curves of spiral galaxies provide convincing evidence that the dark component of matter 'outweighs' the luminous component by a factor of at least 3-10. Particle physics has been very generous (perhaps too generous!) in providing candidate particle species (including, massive neutrinos, sneutrinos, monopoles, photinos, gravitinos, axions, pyrgons, etc.) whose relic abundance (from the big bang) would provide the mass density known to be contributed by the dark component. As I will discuss, the nature of the dark matter bears significantly on the question of galaxy formation and the formation of large-scale structure (voids, superclusters, filaments, etc.) in the Universe. Finally, I will review how astrophysical and cosmological observations have been used to constrain in a significant and very important way the properties (masses, lifetimes, and number of flavors) of neutrinos, the coupling strength of axions (equivalently the Peccei-Quinn symmetry breaking scale), and the flux of relic superheavy magnetic monopoles.

THE STANDARD COSMOLOGY[4]

The hot big bang model nicely accounts for the universal (Hubble) expansion, the 2.7 K cosmic microwave background radiation, and through primordial nucleosynthesis, the abundances of D, ^{4}He and perhaps also ^{3}He and ^{7}Li. Light received from the most distant objects observed (QSOs at redshifts $\approx$ 3.5) left these objects when the Universe was only a few billion years old. Thus observations of galaxies allow us to directly probe the history of the Universe to within a few billion years of 'the bang'. The surface of last scattering for the microwave background is the Universe about 100,000 yrs after the bang when the temperature was about 1/3 eV. Thus the microwave background is a fossil record of the Universe at that very early epoch. In the standard cosmology the epoch of big bang nucleosynthesis takes place from $t \approx 10^{-2}$ sec - 10^2 sec when the temperature was $\approx$ 10 MeV - 0.1 MeV. The light elements synthesized, primarily D, ^{3}He, ^{4}He, and ^{7}Li, are relics from this early epoch, and thus comparing their predicted big bang abundances with their inferred primordial abundances is the most stringent test of the standard cosmology we have at present. [Note that I must say inferred primordial abundance because contemporary astrophysical processes can affect the abundance of these light isotopes, e.g., stars very efficiently burn D, and produce ^{4}He.] At present the predicted abundances of D, ^{3}He, ^{4}He, and ^{7}Li are all simultaneously consistent with their inferred primordial abundances so long as the number of light ($\lesssim$ 1 MeV) neutrino species is less than or equal to 4, and the baryon-to-photon ratio η is in the range[6] (see Figures 1 and 2):

$$\eta = (4\text{-}7) \times 10^{-10}. \qquad (1)$$

The baryon-to-photon ratio is related to the fraction of critical density contributed by baryons by,

$$\eta = 2.83 \times 10^{-8}\, \Omega_b h^2\, (2.7\ \mathrm{K}/T)^3, \qquad (2)$$

where the Hubble parameter $H_0 = 100\ h$ kms^{-1} Mpc^{-1}, and T is the present temperature of the cosmic microwave background. Observations strongly suggest that: $\frac{1}{2} \le h \le 1$ and 2.7 K $\le T \le$ 3.0 K, so that the concordant range for η implies

$$0.014 \le \Omega_b \le 0.15, \qquad (3)$$

i.e., baryons alone cannot provide the closure density.

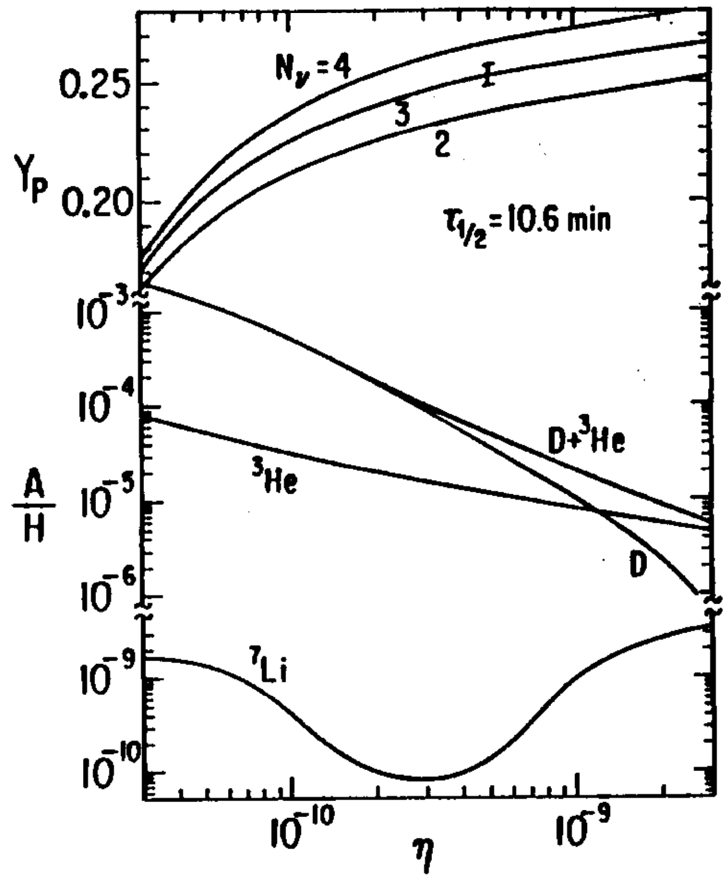


Figure 1 - The predicted primordial abundances of D, ${}^3He$, ${}^4He$, and ${}^7Li$ ($\tau_{1/2}(n) = 10.6$ min was used; error bar shows $\Delta\tau_{1/2} = \pm 0.2$ min; Y_p = mass fraction of ${}^4He$). Inferred primordial abundances: $Y_p = 0.23 - 0.25$; $(D/H)_p \gtrsim 10^{-5}$; $(D + {}^3He)p/H \lesssim 10^{-4}$; $({}^7Li/H)p = (1.1 \pm 0.4) \times 10^{-10}$. Consistency of the predicted abundances with observations can only be achieved for η (= baryon-to-photon ratio) = $(4-7) \times 10^{-10}$ and N_ν (= number of light neutrino species) ≤ 4. For $4 \lesssim \eta/10^{-10} \lesssim 7$, $0.014 \lesssim \Omega_B \lesssim 0.15$. See ref. 6 for more details.

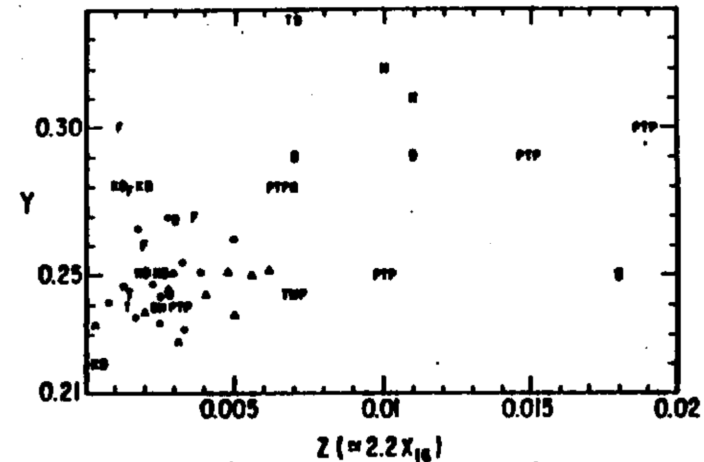
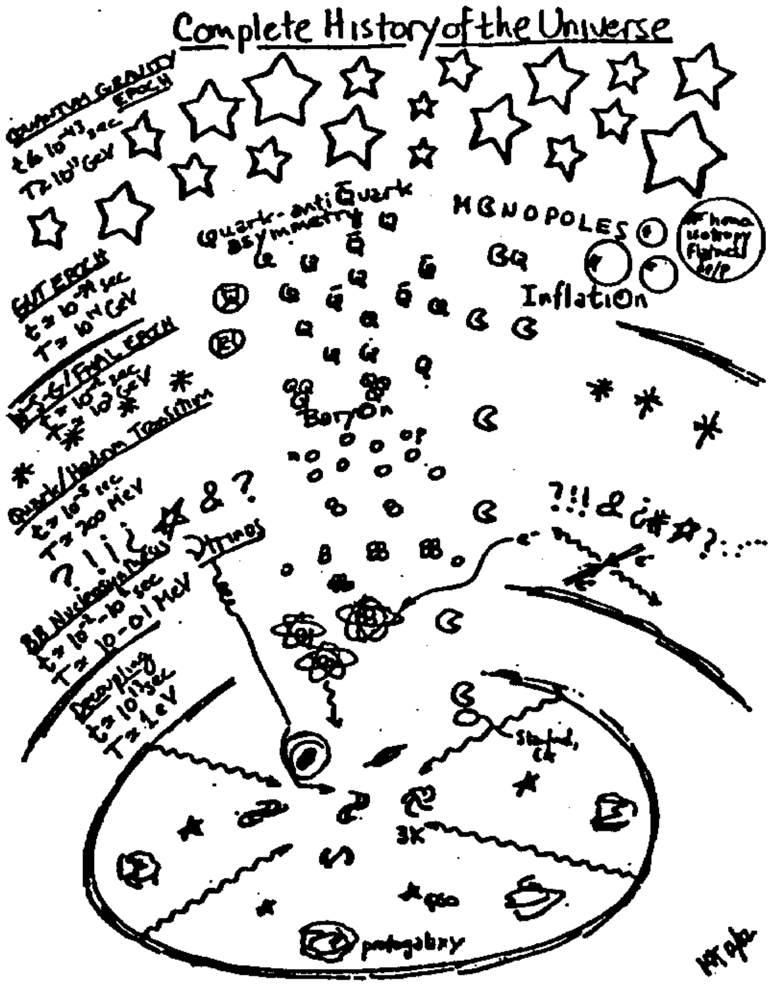


Figure 2 - Summary of determinations of ${}^4He$ mass fraction (Y) in HII regions as a function of the metal abundance Z(--more precisely, 2.2 times the mass fraction of ${}^{16}O$). Where the metal abundance is lowest, one expects the stellar contribution to Y to be the smallest. The data exhibit this trend and clearly show the existence of a primordial ${}^4He$ component of about 0.23-0.25 (by mass). For more details see ref. 6.

Figure 3 - 'The Complete History of the Universe'. Highlights include: decoupling ($t = 10^{13}$ sec, $T = 1/3$eV) - the surface of last scattering for the cosmic microwave background, epoch after which matter and radiation cease to interact and matter 'recombines' into neutral atoms (D, ${}^3He$, ${}^4He$, ${}^7Li$), also marks the beginning of the formation of structure; primordial nucleosynthesis ($t = 10^{-2}$ sec, $T = 10$ MeV) - epoch during which all of the free neutrons and some of the free protons are synthesized into D, ${}^3He$, ${}^4He$, and ${}^7Li$, and the surface of last scattering for the cosmic neutrino backgrounds; quark/hadron transition ($t = 10^{-5}$ sec, T = few 100 MeV)- epoch of 'quark enslavement' (confinement transition in SU(3)); W-S-G epoch ($t = 10^{-12}$ sec, $T = 10^3$ GeV) - SSB phase transition associated with electroweak breaking, $SU(2) \times U(1) \to U(1)$; GUT epoch (??$t = 10^{-34}$ sec, $T = 10^{14}$ GeV??) - SSB of the GUT, during which the baryon asymmetry of the Universe evolves, monopoles are produced, and 'inflation' may occur; the Quantum Gravity Wall ($t = 10^{-43}$ sec, $T = 10^{19}$ GeV).



Note, that other information we have about Ω_b (e.g., lower bound based upon the amount of luminous matter in the Universe, the total amount of matter associated with a galaxy) is consistent with this range. The concordance of the predictions and observations of D and ^{4}He are particularly compelling evidence because there is no known contemporary astrophysical site where the observed amounts amounts of ^{4}He ($\approx$ 25% by mass) and D (D/H $\approx$ few x 10^{-5}) can be produced. It is the successful predictions of big bang nucleosynthesis that gives us confidence in the standard model back to $\approx 10^{-2}$ sec after 'the bang'.

On the large scale (>> 100 Mpc), the Universe is isotropic and homogeneous, and so it can accurately be described by the Robertson-Walker line element[4]

$$ds^2 = -dt^2 + R(t)^2[dr^2/(1-kr^2) + r^2d\theta^2 + r^2\sin^2\theta d\phi^2], \quad (4)$$

where ds^2 is the proper separation between two events, k=1, 0, or -1 is the curvature signature, and R(t) is the cosmic scale factor. The expansion of the Universe is embodied in R(t)--as R(t) increases all proper (i.e., measured by meter sticks) distances scale with R(t), e.g., the distance between two galaxies comoving with the expansion (i.e., fixed r, θ, ϕ), the wavelength of a photon. The k>0 spacetime has positive spatial curvature and is finite in extent; the k<0 spacetime has negative spatial curvature and is infinite in extent; the k=0 spacetime is spatially flat and is also infinite in extent.

The evolution of the cosmic scale factor is determined by the Friedmann equations:

$$H^2 \equiv (\dot{R}/R)^2 = 8\pi G\rho/3 - k/R^2, \quad (5a)$$

$$d(\rho R^3) = -p\, d(R^3), \quad (5b)$$

where ρ is the total energy density and p is the pressure. The expansion rate H (also called the Hubble parameter) sets the characteristic time for the growth of R(t)--H^{-1} = e-folding time for R. The present value of H is 100 h $\mathrm{kms^{-1}Mpc^{-1}}$ = h $(10^{10}$ yr$)^{-1}$ ($\frac{1}{2} \le h \le 1$). As can be seen from eqn. (5a) model Universes with $k \le 0$ expand forever, while a model Universe with $k > 0$ must eventually recollapse. The sign of k (and hence the geometry of spacetime) can be determined from measurements of ρ and H:

$$k/H^2R^2 = \rho/(3H^2/8\pi G) - 1, \quad (6)$$

$$\equiv \Omega - 1,$$

where $\Omega = \rho/\rho_{crit}$ and $\rho_{crit} = 3H^2/8\pi G = 1.88h^2\times10^{-29}\mathrm{gcm^{-3}}$.

From primordial nucleosynthesis we know that $\Omega > \Omega_b \gtrsim 0.014$. The best upper limit to Ω follows by considering the age of the Universe:

$$t_U = 10^{10}\ \mathrm{yr}\ (h^{-1} f(\Omega)), \tag{7}$$

where $f(\Omega) \leq 1$ and is monotonically decreasing. The ages of the oldest stars (in globular clusters) strongly suggest that $t_U \gtrsim 10^{10}$ yr; combining this with eqn. (7) implies that: $\Omega f^2(\Omega) \gtrsim \Omega h^2$. The function Ωf^2 is monotonically increasing and asymptotically approaches $(\pi/2)^2$, implying that independent of h, $\Omega h^2 \lesssim 2.5$. Restricting h to the interval ($\frac{1}{2}$, 1) it follows that: $\Omega h^2 \lesssim 0.8$ and $\Omega \lesssim 3.2$.

The energy density contributed by nonrelativistic matter varies as $R(t)^{-3}$--due to the fact that the number density of particles is diluted by the increase in the proper (or physical) volume of the Universe as it expands. For a relativistic species the energy density varies as $R(t)^{-4}$, the extra factor of R due to the redshifting of the particle's momentum (recall $\lambda \propto R(t)$). The energy density contributed by a relativistic species ($T >> m$) at temperature T is

$$\rho = g_{eff}\pi^2 T^4/30, \tag{8}$$

where g_{eff} is the number of degrees of freedom for a bosonic species, and 7/8 that number for a fermionic species. Note that $T \propto R(t)^{-1}$. Here and throughout I have taken $\hbar = c = k_B = 1$, so that 1 GeV = $(1.97\times10^{-14}\ \mathrm{cm})^{-1}$ = $(1.16\times10^{13}\ \mathrm{K})$ = $(6.57\times10^{-25}\mathrm{sec})^{-1}$, $G = m_{pl}^{-2}$ ($m_{pl} = 1.22\times10^{19}$ GeV), and 1 GeV^4 = 2.32×10^{17} g cm^{-3}.

Today, the energy density contributed by relativistic particles (photons and 3 neutrino species) is negligible: $\Omega_{rel} \approx 4\times10^{-5}h^{-2}(T/2.7K)^4$. However, since $\rho_{rel} \propto R^{-4}$, while $\rho_{nonrel} \propto R^{-3}$, early on relativistic species will dominate the energy density. For $R/R_{today} \lesssim 4\times10^{-5}\ (\Omega h^2)^{-1}(T/2.7K)^4$, which corresponds to $t \lesssim 4\times10^{10}\mathrm{sec}(\Omega h^2)^{-2}(T/2.7K)^6$ and $T \gtrsim 6$ eV $(\Omega h^2)(2.7K/T)^3$, the energy density of the Universe was dominated by relativistic particles. Since the curvature term varies as $R(t)^{-2}$, it too will be small compared to the energy density contributed by relativistic particles, and eqn. (5a) simplifies to: $(\dot{R}/R)^2 = 4\pi^3 g_* T^4/45 m_{pl}^2$, which implies:

$$R(t) \propto t^{\frac{1}{2}} \tag{9a}$$

$$T(t) = 1.6 g_*^{-1/4}(t/10^{-6}\mathrm{sec})^{-\frac{1}{2}}\ \mathrm{GeV}, \tag{9b}$$

(valid for $t \lesssim 10^{10}$ sec, $T \gtrsim 10$ eV). Here g_* counts the total number of effective relativistic degrees of freedom: freedom: $g_*(3K) = 3.36$ ($\gamma, \nu\bar{\nu}$); g_* (few MeV) = 10.75 (γ, $e^{\pm}$, $\nu\bar{\nu}$); g_* (few 100 GeV) ≈ 110 (γ, $W^{\pm}Z$, 8 gluons, 3 families of quarks and leptons, 1 Higgs doublet).

Although our verifiable knowledge of the early history of the Universe only takes us back to $t \approx 10^{-2}$ sec, and $T \approx 10$ MeV, nothing in our present understanding of the laws of physics suggests that it is unreasonable to extrapolate back to times as early as $\approx 10^{-43}$ sec and temperatures as high as $\approx 10^{19}$ GeV. At high energies the interactions of quarks and leptons are asymptotically free (and/or weak) justifying the dilute gas approximation made in eqn. (8), and at energies below 10^{19} GeV quantum corrections to general relativity are expected to be small. I hardly need to remind the reader that 'not unreasonable' does not necessarily mean 'correct'. Making this extrapolation, I have summarized 'The Complete History of the Universe' in Figure 3.

THE MARRIAGE OF GUTs AND THE STANDARD COSMOLOGY

Grand Unified Theories make two startling predictions: (1) interactions which violate baryon and lepton number conservation; (2) the existence of stable, superheavy magnetic monopoles[7]. Both have very important cosmological consequences. The presence of baryon non-conservation in GUTs means that the baryon number of the Universe is not fixed--this turns out to be a very good thing for the standard cosmology. The presence of superheavy magnetic monopoles in GUTs means that there could be relic monopoles left over from a very early epoch--this turns out to be a very bad thing for the standard cosmology.

Baryogenesis - Although the laws of physics are very nearly matter-antimatter symmetric (the only observed violation being that in the K^0-$\bar{K}^0$ system), the Universe apparently is not! There is no evidence for appreciable quantities of antimatter in the Universe: NASA has yet to lose a space probe due to it annihilating with 'anti stuff' in the solar system. Cosmic rays circulate throughout the galaxy, and provide us with samples of material from all parts of the galaxy. The ratio of antiprotons to protons in the cosmic rays is $\approx 10^{-4}$-- although this ratio is not presently consistent with all of the antiprotons being produced in energetic cosmic ray proton collisions, it is consistent with the absence of appreciable amounts of antimatter in the galaxy (for further discussion of the antiproton puzzle see ref. 8). Finally, if clusters of galaxies contained both matter galaxies and antimatter galaxies, then the intracluster gas would contain both matter and antimatter, and would be aglow with γ rays from matter-antimatter annihilations. For nearby clusters like Virgo, these γ rays are not seen. [The issue of antimatter in the Universe is thoroughly reviewed in ref. 9.]

These observations lead us to believe that there is no appreciable quantity of antimatter in the Universe; more precisely, that if there is, it must be separated on scales $\gtrsim 1\ M_\odot$, $10^{12} M_\odot$, 10^{14}-$10^{15} M_\odot$ respectively--this observation alone is essentially impossible to reconcile with any baryon symmetric cosmology. I will assume that $n_b >> n_{\bar{b}}$; then the baryon-to-photon ratio η is also the baryon number ($= n_b - n_{\bar{b}}$)-to-photon number ratio. The number of photons in the Universe has not been conserved, but has increased as various massive particle species annihilated (e.g., $e^\pm$ pairs at $T \approx \frac{1}{2}$ MeV). However, the total entropy has remained constant (assuming that the expansion has been adiabatic). The entropy is dominated by the contribution from relativistic particles (and is $\propto$ number of relativistic particles). The entropy density $s \simeq (4/3)\rho_{rel}/T = (2\pi^2/45) g_* T^3$, and today $s \simeq 7\ n_\gamma$. Thus the baryon number-to-entropy is

$$n_B/s \simeq (1/7)\eta \simeq (6\text{-}10) \times 10^{-11} \qquad (10)$$

--a ratio which remains fixed so long as B is effectively conserved and the expansion is adiabatic. This ratio can be rightly called the baryon number of the Universe (in spite of its rather curious value $\simeq \exp(-\pi^2)$).

Although the matter-antimatter asymmetry appears to be maximal today, $n_B/s \simeq (6\text{-}10) \times 10^{-11}$ implies that earlier on it was very tiny. To see this, let us assume for simplicity that nucleons are the fundamental baryons. Earlier than 10^{-6} sec after 'the bang' the temperature was greater than the mass of a nucleon. Thus nucleons and antinucleons should have been about as abundant as photons, $n_N \simeq n_{\bar{N}} \simeq n_\gamma$. The entropy density s is $\simeq g_* n_\gamma \simeq g_* n_N = O(10^2)\ n_N$. The constancy of $n_B/s \simeq O(10^{-10})$ requires that for $t < 10^{-6}$ sec, $(n_N - n_{\bar{N}})/n_N$ ($\simeq 10^2\ n_B/s$) $\simeq O(10^{-8})$. During its earliest epoch, the Universe was very nearly (but not quite) baryon symmetric.

The non-zero value of the baryon number of the Universe is absolutely essential (for our existence!). Consider a model Universe which is baryon symmetric ($n_B/s = 0$). Earlier than 10^{-6} sec after 'the bang' nucleons and antinucleons were about as abundant as photons. For $T < 1$ GeV the equilibrium abundance of nucleons and antinucleons is $(n_N/n_\gamma)_{EQ} \simeq (m_N/T)^3$ $\exp(-m_N/T)$, and as the Universe cooled the number of nucleons and antinucleons would decrease, tracking the equilibrium abundance as long as the annihilation rate

The necessity of a departure from thermal equilibrium is a bit more subtle. It can be shown that CPT invariance alone guarantees that in thermal equilibrium the number of baryons and antibaryons are equal (even in the presence of interactions which violate B, C, and CP). The Universe is only in thermal equilibrium when the rates for the reactions which drive it to equilibrium are rapid compared to the rate of change of the temperature ($\dot{T}/T = -H$). Departures from equilibrium have occurred frequently (thank God!); e.g., in spite of the low temperature, matter in the Universe is not all in the form of Fe--the most tightly bound nucleus.

The basic idea of baryogenesis has been discussed by many authors; the model which incorporates Sakharov's three ingredients and which has become the 'standard scenario' is the so-called out-of-equilibrium decay scenario. I will now describe the scenario in some detail, though only qualitatively.

Denote by 'S' a superheavy ($\gtrsim 10^{14}$ GeV) boson whose interactions violate B conservation. S might be a gauge or a Higgs boson. Let its coupling strength to fermions be $\alpha^{1/2}$, and its mass be M. From dimensional considerations its decay rate $\Gamma_D = \tau^{-1}$ should be $\Gamma_D \simeq \alpha M$.

At the Planck time ($\simeq 10^{-43}$ sec) assume that the Universe is baryon symmetric ($n_B/s=0$), with all fundamental particle species (fermions, gauge and Higgs bosons) present with equilibrium distributions. At this epoch $T \sim g_*^{-1/4}\ m_{pl} \sim 3 \times 10^{18}$ GeV $>> M$ (here I have taken $g_* \sim O(100)$; in minimal SU(5) $g_* \simeq 160$), so S, $\bar{S}$ bosons are very relativistic and up to statistical factors as abundant as photons: $n_S \simeq n_{\bar{S}} \simeq n_\gamma$. Nothing of importance occurs until $T \simeq M$.

For $T < M$ the equilibrium abundance of S, $\bar{S}$ bosons relative to photons is $n_S/n_\gamma \simeq (M/T)^{3/2} \exp(-M/T)$. [$n_S/n_\gamma$ is just the number of S, $\bar{S}$ bosons per comoving volume.] In order for S, $\bar{S}$ bosons to maintain an equilibrium abundance as T falls below M, they must be able to diminish in number rapidly compared to $H = |\dot{T}/T|$. The most important process in this regard is decay; other processes (e.g., annihilation) are higher order in α. If $\Gamma_D >> H$ for $T \simeq M$, then S, $\bar{S}$ bosons can adjust their abundance (by decay) rapidly enough so that n_S/n_γ 'tracks' its equilibrium value. In this case thermal equilibrium is maintained and no asymmetry is expected to evolve.

More interesting is the case where $\Gamma_D < H \simeq 1.66 g_*^{1/2} T^2/m_{pl}$ when $T \simeq M$, or equivalently $M > g_*^{-1/2} \alpha 10^{19}$ GeV. In this case S, $\bar{S}$ bosons are <u>not</u> decaying on the expansion timescale ($\tau > t$), and so remain as abundant as photons for $T \lesssim M$, and hence are overabundant relative to their equilibrium number. This overabundance is the departure from thermal equilibrium. Much later when $T << M$, $\Gamma_D \approx H$ (i.e., $t \approx \tau$), and S, $\bar{S}$ bosons begin to decrease in number due to decays. To a good approximation they decay freely since the fraction of fermion pairs with sufficient CM energy to produce an S, $\bar{S}$ pair is $\sim \exp(-M/T) << 1$, greatly suppressing inverse decay processes ($\Gamma_{ID} \sim \exp(-M/T)\, \Gamma_D << H$).

Now consider the decay of S and $\bar{S}$ bosons; suppose S decays to channels 1 and 2 with baryon numbers B_1 and B_2, and branching ratios r and (1-r). Denote the corresponding quantities for $\bar{S}$ by $-B_1$, $-B_2$, $\bar{r}$, and $(1-\bar{r})$. [e.g., 1 = ($\bar{q}\bar{q}$), 2 = (ql), $B_1 = -2/3$, and $B_2 = 1/3$.]. The mean net baryon number of the decay products of the S and $\bar{S}$ are respectively:

$$B_S = rB_1 + (1-r)B_2\ ,\ B_{\bar{S}} = -\bar{r}B_1-(1-\bar{r})B_2\ .$$

Hence the decay of an S, $\bar{S}$ pair on average produces a baryon number ϵ,

$$\epsilon \equiv B_S + B_{\bar{S}} = (r-\bar{r})\ (B_1 - B_2)\ .$$

When the S, $\bar{S}$ bosons decay ($T << M$, $t \approx \tau$) $n_S \approx n_{\bar{S}} \approx n_\gamma$. Therefore, the net baryon number density produced is $n_B \approx \epsilon n_\gamma$. The entropy density $s \approx g_* n_\gamma$, and so the baryon asymmetry produced is $n_B/s \approx \epsilon/g_* \approx 10^{-2}\,\epsilon$. This simple picture is borne out by detailed numerical calculations (which among other things indicate that the mass condition for out-of-equilibrium decay is not as stringent as $M > g_*^{-1/2}\alpha\ 10^{19}$ GeV). [For a more detailed discussion of baryogenesis, and a complete list of references, see, e.g., ref. 11.]

The crucial quantity is the C, CP violation in the superheavy system, quantified by $(r-\bar{r})$. Lacking the GUT, $(r-\bar{r})$ and n_B/s cannot be calculated precisely (e.g., as the ^{4}He abundance can). However, some general remarks can be made. It seems very unlikely that $(r-\bar{r})$ can be related to the $K^o - \bar{K}^o$ system (in sign or magnitude), the difficulty being that not enough C, CP violation can be 'fed' up to the superheavy system. In the minimal SU(5) model (3 families, one $\underline{5}$ and one $\underline{24}$ of Higgs) $(r-\bar{r})$ is $<< 10^{-8}$. However, in simple extensions of minimal SU(5) (additional Higgs representations or a fourth family) $(r-\bar{r})$ can easily be $\approx O(10^{-8})$. In more complicated GUTs (e.g., SO(10), E6, etc.) it is also easy to have $(r-\bar{r}) \approx 10^{-8}$, although this requirement restricts possible symmetry breaking patterns. Baryogenesis is not adversely affected by a long proton lifetime (unless $\tau_p \approx \infty$!), and tends to be more difficult in supersymmetric models because of the increased importance of $2 \leftrightarrow 2$ B-nonconserving scattering processes which can erase a baryon asymmetry.

Although baryogenesis is nowhere near being on the same firm footing as primordial nucleosynthesis, we now at least have a very attractive framework for understanding the origin of $n_B/s \approx 10^{-10}$. A framework which is so attractive, that in the absence of observed proton decay, the baryon asymmetry of the Universe may be the best evidence for some kind of quark/lepton unification.

Superheavy Monopoles - In 1974 't Hooft and Polyakov showed that monopoles are obligatory in the low-energy theory whenever a semi-simple group G (e.g., SU(5)) breaks down to a group G'x U(1) which contains a U(1) factor (e.g., SU(3) x SU(2) x U(1)); this, of course, is the goal of unification. These monopoles are associated with nontrivial topology in the Higgs field responsible for SSB, topological knots, if you will, have a mass $O(M/\alpha)$ ($\approx 10^{16}$ GeV in SU(5); $M \approx$ scale of SSB), and magnetic charge which is a multiple of the Dirac charge.

Since there exist no contemporary sites for producing particles of mass even approaching 10^{16} GeV, the only plausible production site is the early Universe, about 10^{-34} sec after 'the bang' when the temperature was $O(10^{14}$ GeV). There are two ways in which monopoles can be produced: (1) as topological defects during the SSB of the unified group G; (2) in monopole-antimonopole pairs by energetic particle collisions. The first process has been studied by Kibble[12], Preskill[13], and Zel'dovich and Khlopov[14], and I will briefly review their important conclusions here. The magnitude of the Higgs field responsible for the SSB of the unified group G is determined by the minimization of the free energy. However, this does not uniquely specify the direction of the Higgs field in group space. A monopole corresponds to a configuration in which the direction of the Higgs field in group space at different points in physical space is topologically distinct from the configuration in which the Higgs field points in the same direction (in group space) everywhere in physical space (which corresponds to no monopole):

no monopole　　　monopole　　　→ = direction of Higgs field in group space.

In the standard, hot big bang cosmology there are particle horizons, i.e., the distance over which a light signal could have propagated since 'the bang' (t=0) is finite and $\approx O(ct)$. At the time of SSB (in SU(5) $T \approx 10^{14}$ GeV and $t \approx 10^{-34}$ sec) the Higgs field can only smoothly orient itself on scales $\lesssim$ horizon $\approx ct \approx 10^{-23}$ cm. This results in O(1) monopole (topological defect) per horizon volume. The horizon volume contains a net baryon number of about $(10^{15}\ \mathrm{GeV}/T)^3$ (corresponding today to that number of baryons). For SU(5) the number of monopoles produced as topological defects is $O(10^{-3})$ per baryon--a seemingly small number when one considers today there are $O(10^{10})$ microwave photons per baryon. Because of ther relative scarcity, monopoles and antimonopoles do not annihilate in appreciable numbers, and for the standard cosmology and SU(5) the relic monopole abundance predicted is $O(10^{-3})$ per baryon. This corresponds to a present mass density of about 10^{12} x the critical mass density! This is clearly impossible. [More precisely, when such a Universe cooled to a temperature of O(3K) it would only be O(30,000 yrs) old!]. This catastrophe is known as the monopole problem. A number of possible remedies have been suggested; all involve either modifying the standard cosmology or the standard GUTs in a nontrivial way. [For a recent review of the monopole problem and possible solutions see ref. 15.] The lesson is clear: the standard cosmology and the simplest GUTs are not compatible. In the next section I will discuss the solution which at present seems to be the most attractive--the new inflationary Universe scenario.

I mention in passing that if the glut of monopoles produced as topological defects in the standard cosmology can be avoided, then the only production mechanism is pair production in very energetic particle collisions, e.g., particle(s) + antiparticle(s) → monopole + antimonopole. [Of course, the 'Kibble production' of monopoles might be consistent with the standard cosmology (and other limits to the monopole flux) if the SSB transition occurred at a low enough temperature, say ≲ $O(10^{10}\ \mathrm{GeV})$.] The numbers produced are intrinsically small because monopole configurations do not exist in the theory until SSB occurs ($T_c \approx M$ = scale of SSB), and have a mass $O(M/\alpha) \approx 100\ M \approx 100\ T_c$. For this reason they are never present in equilibrium numbers; however, some are produced due to the rare collisions of particles with sufficient energy.[15,16] This results in a present monopole to photon ratio

$$n_M/n_\gamma \approx 10^3\ (m/T_{max})^3\ \exp(-2m/T_{max}), \qquad (11)$$

where m is the mass of the monopole, and T_{max} is the highest temperature achieved after SSB. For reference, $\Omega_M \approx 3 \times 10^{23}\ (n_M/n_\gamma)(m/10^{16}\ \mathrm{GeV})$, and the monopole flux $F \approx 10^9 (n_M/n_\gamma)\ \mathrm{cm}^{-2}\mathrm{sr}^{-1}\mathrm{s}^{-1}$. In general $m/T_{max} \approx O(100)$ so that $n_M/n_\gamma \approx O(10^{-70})$--a negligible number of monopoles. However, the number produced is exponentially sensitive to m/T_{max}, so that a factor of 3-10 uncertainty introduces an enormous uncertainty in the predicted production (for further discussion of this point see ref. 17).

Cosmology seems to leave the poor monopole hunter/huntress with two firm predictions: that there should be equal numbers of north and south poles; and that either far too few to detect, or far too many to be consistent with the standard cosmology should have been produced. The detection of any superheavy monopoles would necessarily send theorists back to their chalkboards!

Two messages emerge very clearly from the 'marriage of the standard cosmology with GUTs': B nonconservation is very good for cosmology; and the simplest GUTs and the standard cosmology are incompatible--meaning that one (or both) have to be modified. Both messages are very significant bits of information about physics at very high energy and the earliest moments of the Universe.

THE INFLATIONARY UNIVERSE SCENARIO

The hot big bang cosmology is, by any standard, a remarkable achievement. However, it has its shortcomings. By itself, it fails to account for or even elucidate a number of fundamental cosmological facts. They include: the large-scale homogeneity and isotropy of the Universe, the origin of the small-scale inhomogeneity, the near critical expansion rate of the Universe today, the predominance of matter (GUTs do provide us with a framework for understanding this), 'the monopole problem', and the extreme smallness of the present value of the cosmological term (as measured in any natural system of units). Due in large measure to Guth's inflationary Universe paradigm[1], great progress has been made in recent years toward understanding these 'cosmological facts' at a very fundamental level. I will begin this section by reviewing the cosmological conundrums, and then I'll go on to discuss new inflation[2,3], the variant of Guth's original scenario which seems to be capable of resolving all but one of the puzzles.

Large-Scale Homogeneity and Isotropy - The observable Universe ($d \approx H^{-1} \approx 10^{28}$ cm $\approx$ 3000 Mpc) is to a high degree of precision isotropic and homogeneous on the largest scales (>> 100 Mpc). The best evidence for this is provided by the uniformity of the cosmic background temperature: $\Delta T/T \lesssim 10^{-3}$ (see Fig. 4 below). Large-scale density inhomogeneities or an anisotropic expansion would result in fluctuations in the microwave background temperature of a comparable size.[18,19] The smoothness of the observable Universe is puzzling if one wishes to understand it as a result of microphysical processes operating in the early Universe. As mentioned earlier the standard cosmology has particle horizons, and when matter and radiation last vigorously interacted (decoupling: $t \approx 10^{13}$ s, $T \approx 1$ eV) what was to become the present observable Universe was comprised of $\approx 10^6$ causally-distinct regions. Put slightly differently,

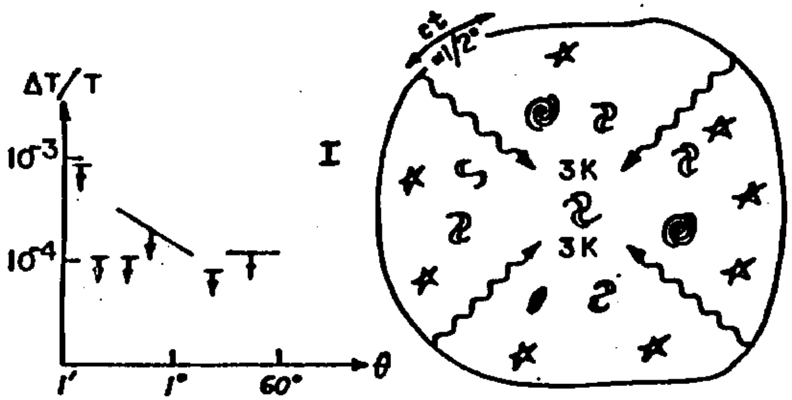


Figure 4 - Summary of the measurements of the fluctuations in the microwave background temperature on angular scales $\gtrsim 1'$ (from D. Wilkinson, 1983).

the particle horizon at decoupling only subtends an angle of about $\frac{1}{2}°$ on the sky today; how is it that the microwave background temperature is so uniform on angular scales $\gg \frac{1}{2}°$?

Small-Scale Inhomogeneity - As any astronomer will gladly tell you, on small scales ($\lesssim$ 100Mpc) the Universe is very lumpy (stars, galaxies, clusters of galaxies, etc.). The uniformity of the microwave background on very small angular scales ($\ll 1°$) indicates that the Universe was smooth, even on these scales at the time of decoupling. [Note, today $\delta\rho/\rho \approx 10^5$ on the scale of a galaxy.] Whence came the structure that is so conspicuous today? Once matter decouples from the radiation and is free of the pressure support provided by the radiation, small inhomogeneities will grow via the Jeans (gravitational) instability: $\delta\rho/\rho \propto t^{2/3} \propto R$ (in the linear regime). [If the mass density of the Universe is dominated by a collisionless particle species, e.g.,

a light relic neutrino species, or axions, density perturbations in these particles can begin to grow when the Universe becomes matter-dominated, $R \approx 4 \times 10^{-5}\ R_{today}$ for $\Omega h^2 = 1$.] Density perturbations of amplitude $\delta\rho/\rho \approx 10^{-3}$ or so, on the scale of a galaxy ($\approx 10^{12} M_\odot$) at the time of decoupling seem to be required to account for the small-scale structure observed today. Their origin, their spectrum (certainly perturbations should exist on scales other than $10^{12} M_\odot$), their nature (adiabatic or isothermal), and the composition of the dark matter are all crucial questions for understanding the formation of structure, which to date remain unanswered.

Flatness/Near Critical Expansion Rate - As I discussed earlier Ω is known to be in the range: 0.014-3.2. Using eqn. (5a), Ω can be written as

$$\Omega = 1/(1-x(t)), \qquad (12a)$$

$$x(t) = (k/R^2)/(8\pi G\rho/3), \qquad (12b)$$

and the expansion rate as

$$H = H_{crit}(1-x(t))^{\frac{1}{2}}, \qquad (13a)$$

$$H_{crit}^2 = 8\pi G\rho/3. \qquad (13b)$$

[H_{crit} is the expansion rate in the k=0 model--a model Universe which expands forever, albeit at an ever decreasing rate.] Note that neither Ω nor H/H_{crit} are constant; they both vary with time since $x(t) \propto R(t)^n$ (n=1 - matter-dominated; n=2 - radiation-dominated). Since $\Omega \approx O(1)$ today, x_{today} can be at most O(1). This implies that at the epoch of nucleosynthesis: $x_{BBN} \lesssim 10^{-16}$, $\Omega_{BBN} = 1 \pm O(\lesssim 10^{-16})$, and $(H/H_{crit})_{BBN} = 1 \pm O(\lesssim 10^{-16})$, and extrapolating back to the Planck epoch that: $x_{Pl} \lesssim O(10^{-60})$, $\Omega_{Pl} = 1 \pm O(\lesssim 10^{-60})$, and $(H/H_{crit})_{Pl} = 1 \pm O(\lesssim 10^{-60})$. That is, very early on the ratio of the curvature term to the density was extremely tiny, or equivalently, the expansion of the Universe proceeded at the critical rate to a high degree of precision. Since x(t) has apparently always been $\lesssim 1$, our Universe is today and has been in the past closely-modeled by the k=0 flat cosmology. Were the ratio x not exceedingly small early on, the Universe would have either recollapsed long ago (for k > 0), or began a 'coasting phase' (for k < 0) where $R \propto t$. [If k < 0 and $x_{Pl} = 1$, then T = 3K for $t \approx 10^{-11}$ sec!] The smallness of the ratio x required as 'initial data' for our Universe is puzzling (to say the least).

<u>Predominance of Matter/The Monopole Problem</u> - GUTs go a long way toward 'explaining' the predominance of matter in the Universe; however, 'the price', in the context of the standard cosmology, is the glut of monopoles produced at the same epoch.

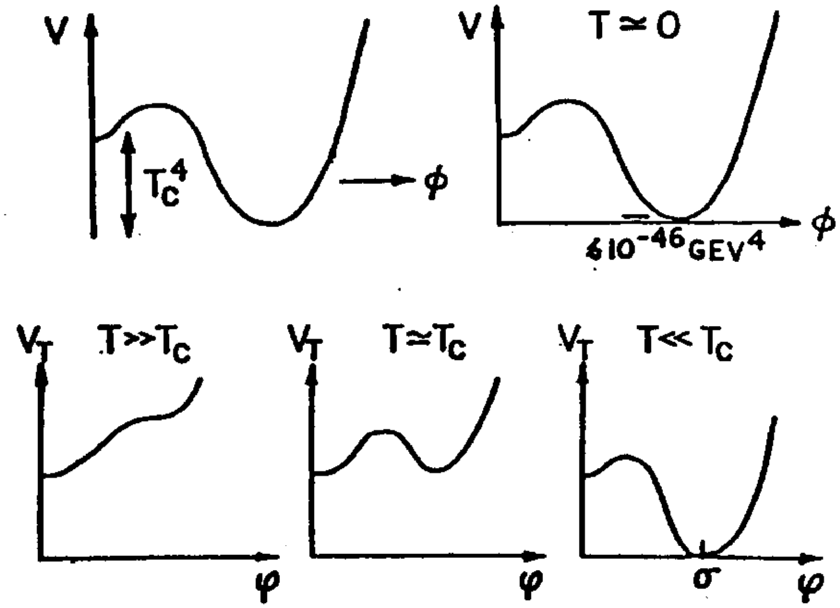
<u>The Cosmological Constant</u> - With the possible exception of supersymmetry and supergravity theories, the absolute scale of the effective potential $V(\phi)$ is not determined in gauge theories (ϕ = one or more Higgs field). At low temperatures $V(\phi)$ is equivalent to a cosmological term (<u>i.e.</u>, contributes $Vg_{\mu\nu}$ to the stress energy of the Universe). The observed expansion rate of the Universe today ($H = 50 - 100$ km s^{-1} Mpc^{-1}) limits the total energy density of the Universe to be $\lesssim O(10^{-29} g\ cm^{-3}) \approx 10^{-46}\ GeV^4$. Thus empirically the vacuum energy of our $T = 0$ $SU(3) \times U(1)$ vacuum ($= V(\phi)$ at the SSB minimum) must be $\lesssim 10^{-46}\ GeV^4$. Compare this to the difference in energy density between the false ($\phi = 0$) and true vacua, which is $O(T_c^4)$ (T_c = symmetry restoration temperature): for $T_c \approx 10^{14}$ GeV, $V_{SSB}/V(\phi = 0) \lesssim 10^{-102}$! At present there is no satisfactory explanation for the vanishingly-small value of the $T = 0$ vacuum energy density (equivalently, the cosmological term).

Today, the vacuum energy is apparently negligibly small and seems to play no significant role in the dynamics of the expansion of the Universe. If we accept this empirical determination of the absolute scale of $V(\phi)$, then it follows that the energy of the false ($\phi = 0$) vacuum is enormous ($\approx T_c^4$), and thus could have played a significant role in determining the dynamics of the expansion of the Universe. Accepting this very non-trivial assumption about the zero of the vacuum energy is the starting point for inflation (see Fig. 5).

<u>Generic New Inflation</u> - The basic idea of the inflationary Universe scenario is that there was an epoch when vacuum energy density dominated the energy density of the Universe. During this epoch $\rho \approx V \approx$ constant, and thus $R(t)$ grows exponentially ($\propto \exp(Ht)$), allowing a small, causally-coherent region (initial size $\lesssim H^{-1}$) to grow to a size which encompasses the region which eventually becomes our presently-observable Universe. In Guth's original scenario,[1] this epoch occurred while the Universe was trapped in the false ($\phi = 0$) vacuum during a strongly first-order phase transition. Unfortunately, in models which inflated enough (<u>i.e.</u>, underwent sufficient exponential expansion) the Universe never made a 'graceful return' to the usual radiation-dominated FRW cosmology[20]. Rather than discussing the original model and its shortcomings in detail, I will instead focus on the variant, dubbed 'new inflation', proposed independently by Linde[2] and Albrecht and Steinhardt[3]. In this scenario, the vacuum-dominated

epoch occurs while the region of the U[illegible]e in question is slowly, but inevitably, evolving t[illegible]he true, SSB vacuum. Rather than considering specific models in this section, I will discuss new inflation for a generic model.

Consider a SSB phase transition which occurs at an energy scale M_G. For $T \gtrsim T_c \approx M_G$ the symmetric ($\phi = 0$) vacuum is favored, <u>i.e.</u>, $\phi = 0$ is the global minimum of the finite temperature effective potential $V_T(\phi)$ (= free energy density). As T approaches T_c a second minimum develops at $\phi \neq 0$, and at $T = T_c$ the two minima are degenerate. [I am assuming that this SSB transition is a first-order phase transition.] At temperatures below T_c the SSB ($\phi = \sigma$) minimum is the global minimum of $V_T(\phi)$ (see Fig. 5). However, the Universe does not instantly m[illegible]he transition from $\phi = 0$ to $\phi = \sigma$; the details and



<u>Figure 5</u> - The finite temperature effective potential at various temperatures (T_c = the critical temperature for the SSB transition).

time required are a question of dynamics.

Assuming a barrier exists between the false and true vacua, thermal fluctuations and/or quantum tunneling or the loss of metastability, must be responsible for taking ϕ across the barrier. The dynamics of this process determine when and how the process occurs (bubble formation, spinodal decomposition, etc.) and the value of ϕ after the barrier is penetrated. For definiteness suppose that the barrier is overcome when the temperature is T_{MS} (which could in principle be $\approx M_G$) and the value of ϕ is ϕ_0. From this point the journey to the true vacuum is downhill (literally) and the evolution of ϕ should be adequately described by the semi-classical equations of motion for ϕ:

$$\ddot{\phi} + 3H\dot{\phi} + \Gamma\dot{\phi} + V' = 0, \qquad (14)$$

where ϕ has been normalized so that its kinetic term in the Lagrangian is $\frac{1}{2}\partial_\mu\phi\partial^\mu\phi$, and prime indicates derivative with respect to ϕ. The subscript T on V has been dropped; for $T \ll T_c$ the temperature dependence of V_T can be neglected and the zero temperature potential ($\equiv V$) can be used. The $3H\dot{\phi}$ term acts like a frictional force, and arises because the expansion of the Universe 'redshifts away' the kinetic energy of $\phi(\propto R^{-3})$. The $\Gamma\dot{\phi}$ term accounts for particle creation due to the time-variation of ϕ (refs. 21, 22). The quantity Γ is determined by the particles which couple to ϕ and the strength with which they couple (Γ^{-1} = lifetime of a ϕ particle). The expansion rate H is determined by the energy density of the Universe [through eqn. (5a)]:

$$\rho = \frac{1}{2}\dot{\phi}^2 + V(\phi) + \rho_r, \qquad (15)$$

where ρ_r represents the energy density in radiation produced by the time variation of ϕ. For $T_{MS} \ll T_c$ the original thermal component makes a negligible contribution to ρ. The evolution of ρ_r is determined by

$$\dot{\rho}_r + 4H\rho_r = \Gamma\dot{\phi}^2, \qquad (16)$$

where the $\Gamma\dot{\phi}^2$ term accounts for particle creation by ϕ.

In writing eqns. (14-16) I have implicitly assumed that ϕ is spatially homogeneous. In some small region (inside a bubble or a fluctuation region) this will be a good approximation. The size of this smooth region will be unimportant; take it to be of order the 'physics horizon', H^{-1}. Now let's follow the evolution of ϕ within the small, smooth patch of size H^{-1}. [I call H^{-1} the 'physics horizon' because coherent physical processes can only take place on a timescale $\lesssim H^{-1}$ = e-folding time for R(t), and thus causality implies a length scale

$\lesssim H^{-1}$ on which coherent physical processes can operate.]

If the potential is flat somewhere between $\phi = \phi_0$ and $\phi = \sigma$, then ϕ will evolve very slowly in that region, and the motion of ϕ will be 'friction-dominated' so that $3H\dot{\phi} = -V'$ (in the slow growth phase particle creation is not important). If V is sufficently flat, then the time required for ϕ to transverse the flat region can be long compared to the expansion timescale H^{-1}, say for definiteness, $\tau_\phi = 100\ H^{-1}$. During this slow growth phase $\rho \approx V(\phi) \approx V(\phi = 0)$; both ρ_r and $\frac{1}{2}\dot{\phi}^2$ are $\ll V(\phi)$. The expansion rate H is then just

$$H = (8\pi V(0)/3m_{pl}^2)^{\frac{1}{2}} \qquad (17)$$
$$\approx M_G^2/m_{pl},$$

where V(0) is assumed to be of order M_G^4. While H $\approx$ constant R grows exponentially: $R \propto \exp(Ht)$; for $\tau_\phi \approx 100\ H^{-1}$ R expands by a factor of e^{100} during the slow rolling period, and the physical size of the smooth region increases to $e^{100}H^{-1}$. This exponential growth phase is called a deSitter phase.

As the potential steepens, the evolution of ϕ quickens. Near $\phi = \sigma$, ϕ oscillates around the SSB minimum with frequency ω: $\omega^2 \approx V''(\sigma) \approx M_G^2 \gg H^2 \approx M_G^4/m_{pl}^2$. As ϕ oscillates about $\phi = \sigma$ its motion is damped by particle creation and the expansion of the Universe. If $\Gamma^{-1} \ll H^{-1}$, the coherent field energy density $(V + \frac{1}{2}\dot{\phi}^2)$ is converted into radiation in less than an expansion time ($\Delta t_{RH} \approx \Gamma^{-1}$), and the patch is reheated to a temperature $T \approx O(M_G)$--the vacuum energy is efficiently converted into radiation ('good reheating'). On the other hand, if $\Gamma^{-1} \gg H^{-1}$, then ϕ continues to oscillate and the coherent field energy redshifts away with the expansion: $(V + \frac{1}{2}\dot{\phi}^2) \propto R^{-3}$. [The coherent field energy behaves like nonrelativistic matter; see ref. 23.] Eventually, when $t \approx \Gamma^{-1}$ the coherent field oscillations decay into light particles, and the patch is reheated to a temperature $T \approx (\Gamma/H)^{\frac{1}{2}}M_G \approx (\Gamma m_{pl})^{\frac{1}{2}} \ll M_G$ ('poor reheating'). The evolution of ϕ is summarized in Fig. 6.

For simplicity, let us assume 'good reheating' ($\Gamma \gg H$). After reheating the patch has a physical size $e^{100}H^{-1}$ ($\approx 10^{17}$ cm for $M_G \approx 10^{14}$ GeV), is at a temperature of order M_G, and in the approximation that ϕ was initially constant throughout the patch, the patch is exactly smooth. From this point forward the region evolves like a radiation-dominated FRW model. How have the cosmological conundrums been 'explained'? First, the homogeneity and isotropy; our observable Universe today ($\approx 10^{28}$ cm) had a physical size of about 10 cm ($\approx 10^{28}$cm$\times$3K/10^{14}GeV) when T was 10^{14} GeV. Thus it lies well within one of the smooth regions produced by

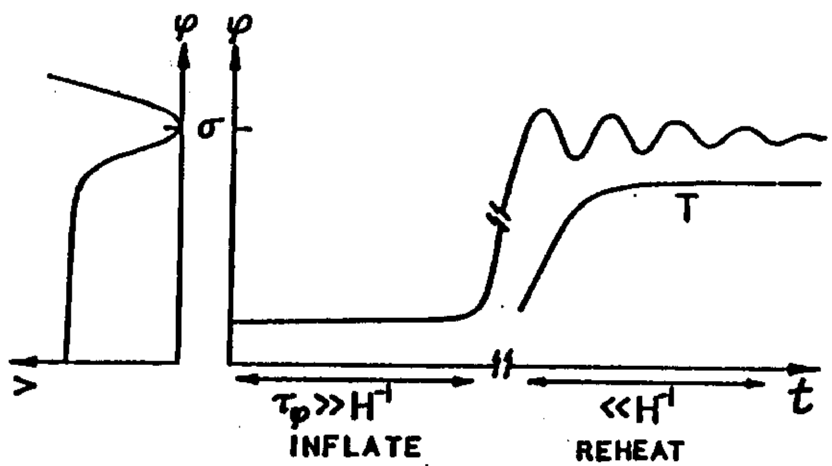


Figure 6 - The evolution of $\phi(t)$. During the slow growth phase the time required for ϕ to change appreciably is $\gg H^{-1}$. As the potential steepens ϕ evolves rapidly (timescale $\ll H^{-1}$), eventually oscillating about the SSB minimum. Particle creation damps the oscillations of ϕ in a time $\approx \Gamma^{-1}$, reheating the Universe to a temperature $\approx M_G$ if $\Gamma \gg H$.

the inflationary epoch. The flatness puzzle involves the smallness of the ratio of the curvature term to the energy density term. This ratio is exponentially smaller after inflation: $x_{after} \approx e^{-200}\, x_{before}$ since the energy density before and after inflation is $O(M_G^4)$, while k/R^2 has decreased exponentially (by e^{200}). Since the ratio x is reset to an exponentially small value, the inflationary scenario predicts that today Ω should be $1 \pm O(10^{-BIG\,\#})$. Assuming the Universe is reheated to a temperature of order M_G, a baryon asymmetry can evolve in the usual way, although the quantitative details may be slightly different.[11] Alternatively, the baryon asymmetry might be produced directly in the decay of the Higgs field oscillations[24]. Of course, it is

absolutely necessary to have baryogenesis occur after reheating since any baryon number (or any c[illegible] quantum number) present before inflation is diluted by a factor $\approx \exp(3H\tau_\phi)$ - the factor by which the total entropy increases.

In fact, the real key to understanding how inflation works is the entropy production. The entropy of the observable Universe is $\approx 10^{88}$ ($\approx$ 7 times the number of photons within 10^{28} cm of us). In the inflationary scenario a smooth patch of size $\approx H^{-1} \approx 10^{-23}$ cm and initial temperature $\approx 10^{14}$ GeV with total entropy $\approx 10^{15}$ grows to a final size $\exp(H\tau_\phi)\, H^{-1}$ ($\approx 10^{17}$ cm in my example) with temperature $\approx$ initial temperature, corresponding to an entropy increase of $\exp(3H\tau_\phi)$ and a final total entropy of $\approx 10^{135} \gg 10^{88} \approx$ the total entropy of the presently observed portion of the Universe.

Since the patch that our observable Universe lies within was once (at the beginning of inflation) causally-coherent, the Higgs field could have been aligned throughout the patch (indeed, this is the lowest energy configuration), and thus there is likely to be $\lesssim 1$ monopole within the entire patch which was produced as a topological defect. The glut of monopoles which occurs in the standard cosmology does not occur. The production of other topological defects (such as domain walls, etc.) is avoided for similar reasons.] Some monopoles will be produced after reheating in rare, very energetic particle collisions. As discussed earlier, the number produced is exponentially small (and exponentially uncertain).

As described, the inflationary scenario produces a large ($\gg 10^{28}$ cm) region which is absolutely uniform. This in itself is a great achievement--previously it was necessary to assume that the Universe was initially isotropic and homogeneous, an assumption which is somewhat disturbing since the set of initial data for the Einstein equations which evolve to a model Universe as smooth as ours is a set of measure zero.[25] But what about the small density inhomogeneities needed for the eventual formation of structure? The exact smoothness of the final patch results from the assumption that the Higgs field ϕ was precisely spatially uniform within the patch. Due to quantum mechanical fluctuations it is not possible for ϕ to be precisely uniform. In fact, it has been shown that the deSitter space produced fluctuations in ϕ ($\Delta\phi \approx H/2\pi$) lead to a spectrum of density inhomogeneities, which eventually cross the horizon with an amplitude

$$(\delta\rho/\rho)_H \approx H^2/\dot{\phi}\,, \qquad (18)$$

where $\dot{\phi}$ is to be evaluated when the scale in question crossed outside the horizon during the inflationary phase.[26-29] Since $\dot{\phi}$ ≈ constant during inflation, the predicted spectrum of density inhomogeneities is approximately independent of the mass scale (this is the so-called Zel'dovich spectrum). This marks the first time that the spectrum of density inhomogeneities could be calculated from first principles--in the past the spectrum of density perturbations has been taken as arbitrary initial data. An amplitude of $O(10^{-4}-10^{-5})$ would lead to a picture of galaxy formation which is consistent with present observations of the distribution of galaxies, etc., and results in microwave temperature fluctuations (on angular scales >> 1°) which are consistent with current measurements/upper limits. [Such a spectrum results in definite predictions for $\delta T/T$ on large-angular scales, predictions which could be checked when experimental sensitivities improve by about a factor of 10, see ref. 19.] The density perturbation constraint, $H^2/\dot{\phi} = O(10^{-4}-10^{-5})$, is a very stringent constraint on models of inflation.

The new inflationary Universe scenario appears to have the potential to resolve all the cosmological puzzles I have discussed, with the very notable exception of the cosmological constant puzzle. [In fact, in some sense it is a house of cards built upon this puzzle.] Inflation results in a Universe whose present state on scales at least as large as 10^{28} cm is insensitive to initial data (data which we are not ever likely to be privy to) and depends only on microphysics. Globally (i.e., on scales $\gtrsim$ the present size of a bubble or fluctuation region) the Universe could appear very different--highly inhomogeneous and anisotropic.

Although the scenario is extremely attractive, at present, a truly compelling inflationary scenario is still lacking. The first model studied, the Coleman-Weinberg SU(5) model,[2,3] is beset by numerous problems, the worst of which is the amplitude of the predicted density perturbations: $(\delta\rho/\rho)_H = O(10-100)$. A number of groups[30-32] have studied supersymmetric and/or supergravity models, in which the inflationary phase transition is not associated with the SSB of the GUT. To date these models have either been plagued by cosmological problems (e.g., insufficient reheating for baryogenesis[30]), or are just toy models constructed to give a viable inflationary cosmology. So at present we have a very attractive paradigm,[1-3] a very precise prescription for the scalar potential which is wanted,[24] but no compelling particle physics model. It remains to be seen if a model will be found which can both implement new inflation and predict sensible particle physics.

DARK MATTER IN THE UNIVERSE

One of the most fundamental cosmological quantities $\langle\rho\rangle$, the average mass density, is also one of the most poorly constrained. As I discussed earlier Ω $(\equiv\langle\rho\rangle/\rho_{crit})$ is only known to be in the interval 0.014-O(3). How does one determine $\langle\rho\rangle$? From a theorist's point-of-view it's all very straightforward:

$$\langle\rho\rangle = \langle n_{GAL}\rangle \times \langle M_{GAL}\rangle, \qquad (19)$$

i.e., one determines the average number density of galaxies and multiplies it times the average mass associated with a galaxy. Note, by doing this I have tacitly assumed that light faithfully traces mass-a highly non-trivial assumption. However, since our knowledge of the Universe is derived primarily from photons, it is difficult to find a technique which does not rely upon this assumption. ['Real astronomers' measure $\langle\rho\rangle$ by determining the average mass-to-light ratio associated with a galaxy (i.e., measure M, measure L, and take the ratio), and multiplying it times the luminosity density --a quantity which they can also readily measure.]

The basic dynamical technique by which the mass of a distant object (from the earth, sun, on up to cosmological objects) is measured relies upon Kepler's third law:

$$GM = r \times v^2, \qquad (20)$$

here r and v are the orbital radius and velocity of a test particle which orbits the mass M. The mass associated with the luminous stuff (undoubtedly baryons) can be determined by studying the orbits of stars and gas clouds at the distance from the center of a galaxy where the light 'craps out' (Holmberg radius ≈ 10-30 kpc for a spiral galaxy; note that the luminosity falls off exponentially with distance from the center). The luminous mass determined this way corresponds to Ω_{lum} ≈ 0.007, which is easily consistent with nucleosynthesis $(0.014 \lesssim \Omega_b \lesssim 0.15)$ and already indicates the presence of dark baryonic matter (at this level, very likely gas and dust).

If luminous mass were the whole story, then as one measured orbital velocities beyond the Holmberg radius, one would expect these velocities to decrease as $r^{-1/2}$. This is not the case. Orbital (also called rotational) velocities for spiral galaxies remain constant beyond the point where the light has 'crapped out'--indicating that the mass (which cannot be due to luminous stuff) continues to increase linearly with r, cf., eqn. (20).

Rotation curves have been measured out to 120 kpc and continue to be flat. The dark matter inferred by these measurements is at least O(3-10) times the luminous component: $\Omega_{halo} \gtrsim 0.03$-$0.10$.

Several comments are in order at this point. Since there is as of yet no convincing evidence for even a single rotation curve which eventually begins to fall, this is a lower limit to the dark component. It is possible that all of this dark matter is baryonic (and very likely that at least some substantial fraction of it is) since primordial nucleosynthesis tells us that: $0.014 \leq \Omega_b \leq 0.15$. While it is generally assumed that this dark halo material is distributed more or less spherically, there is little or no direct evidence to support this hypothesis. The main reason for this belief is that this material which is less condensed than the luminous matter is likely to be composed of effectively collisionless particles (e.g., low mass stars, black holes, an exotic relic particle species) which cannot undergo the dissipation necessary to form a highly non-spherical structure (such as the galactic disk).

There is additional evidence for the existence of dark matter based on the dynamics of larger systems (binary galaxies, small groups of galaxies, and rich clusters of galaxies). In fact, there is some indication that the inferred value of Ω increases with scale (suggesting that dark matter preferentially clusters on larger scales). Unfortunately for the advocates of $\Omega = 1$ (myself included) there is no dynamical evidence for $\Omega = 1$ (the largest inferred values of Ω are $\approx$ 0.2-0.6). What is very clear is that the dominant component of mass density is dark (be it dark baryons or exotic relics). Faber and Gallagher[33] have recently reviewed the issue of dark matter in the Universe.

What is the dark matter? Let me begin with the more mundane candidates and then proceed to the exotic. If $\Omega \lesssim 0.15$, then the dark matter could be ordinary baryons in the form of Jupiters, low mass stars (the luminosity of a star $\propto M^4$), dead stars (e.g., white dwarfs, neutron stars, or black holes), or even large black holes (of mass up to $10^6\ M_\odot$)--although 'hiding' this much mass in any of these forms is not an easy task![34] If $\Omega \gtrsim 0.15$, then it cannot be ordinary baryons, but could be in the form of primordial black holes ($M \lesssim 10^6\ M_\odot$) since they do not contribute to Ω_b.

Particle physics has generously provided us with a very long list of (mostly hypothetical) weakly-interacting particle species whose relic abundance could supply the mass density contributed by dark matter. A partial listing of the candidates is given below in Table I.

TABLE I - CANDIDATE 'DARKONS'

Candidate	Mass	Abundance
axion	$O(10^{-5}\text{eV})$	$O(10^9\text{cm}^{-3})$
neutrino	$O(30\text{ eV})$	100 cm^{-3}
gravitino/photino	O(keV)	few cm^{-3}
sneutrino	O(GeV)	$O(10^{-6}\text{cm}^{-3})$
photino	O(GeV)	$O(10^{-6}\text{cm}^{-3})$
superheavy monopoles	$O(10^{16}\text{GeV})$	$O(10^{-22}\text{cm}^{-3})$
pyrgons, maximons, newtorites, perrypoles,	$\gtrsim O(10^{19}\text{GeV})$	$\lesssim O(10^{-25}\text{cm}^{-3})$
primordial black holes, jupiters, etc.	$\gg 10^{15}\text{g} = 10^{39}$ GeV	$\ll O(10^{-45}\text{ cm}^{-3})$

Given that dark matter dominates luminous matter by a factor of at least 3-10, the composition of the dark matter will clearly have an important bearing on the formation of structure in the Universe. In spite of the long list of candidates with very different masses some general themes have emerged. For collisionless dark matter (which includes all of the candidates under consideration) there are two limiting cases: 'cryons'[35] (or 'cold' dark matter) and 'thermons' (or 'hot' dark matter). The distinction has to do with whether or not the 'darkon' species is relativistic when the Universe becomes matter-dominated. Recall that an $\Omega = 1$ Universe becomes matter-dominated when $R/R_{today} \approx 4\times10^{-5}$ and the temperature of the Universe was O(10 eV) (regardless of the composition of the dark component); the mass contained within the horizon at that time is $\approx 10^{15}\ h^{-4}\ M_\odot$ (about the mass of a supercluster). 'Thermons' are 'darkon' species which are still relativistic at the time of matter domination--since $T \approx O(10\text{ eV})$, this implies that m_d must be $\lesssim O(30\text{ eV})$ (average energy/particle of a particle species in equilibrium $\approx 3T$). 'Cryons' are particles which are non-relativistic at the epoch of matter-domination, implying that m_d must be $\gg O(30\text{ eV})$ in this case. The exception to this simple mass criterion are axions--they are never in thermal equilibrium and are created very cold (a condensate of zero momentum particles). This leaves the massive neutrino as the only 'thermon' in Table I.

The importance of whether or not the dark particles are relativistic has to do with the damping of small-scale perturbations by free-streaming of the darkons, also referred to as collisionless phase mixing or Landau damping. In an expansion time, perturbations in the darkons can diffuse away on length scales $\lesssim H^{-1}$ x (velocity of the particle); since $H^{-1} \propto t$ the damping scale increases until the darkons become nonrelativistic, at which time the damping scale reaches its maximum value (up to log factors). For a darkon species which is in thermal equilibrium, that maximum damping scale is approximately[36,37]

$$\begin{aligned} M_D &= m_{pl}^3/m_d^2 \\ &= 3\times10^{12} M_\odot / m_d^2(\text{keV}). \end{aligned} \qquad (21)$$

All initial density perturbations on scales $\lesssim M_D$ are washed out by 'free-streaming' of the darkons.

For a thermon, $M_D = 3\times10^{15} M_\odot$, implying that the first structures to form must be supercluster size or larger, and that these objects must subsequently fragment to form galaxies.

On the other hand, for a cryon, $M_D \lesssim 10^{12} M_\odot$ (taking as the criterion $m_d \gtrsim$ keV). In a cryon-dominated Universe structure can form on small scales (galaxies or perhaps smaller) first and then build up to larger-scales in a hierarchical manner. Clearly the thermon and cryon scenarios are qualitatively very different.

Numerical simulations of both scenarios have been performed. The thermon (neutrino) simulations do a fine job of reproducing the large-scale structure and segregating the dark and luminous matter (which is not surprising since the thermons are smooth on scales $\lesssim 10^{15} M_\odot$). This is particularly good for '$\Omega = 1$ advocates' since there is good evidence that the amount of matter which clusters on these scales ($\lesssim 10^{15} M_\odot$) comes nowhere near giving $\Omega = 1$ (closer to $\Omega = 0.1$). However, as one might have expected, these scenarios have difficulty reproducing the observed small-scale structure (in particular, the galaxy-galaxy correlation function).[38]

The cryon simulations look very promising. With the Zel'dovich spectrum of density perturbations they seem to be able to reproduce both the large-scale and the small-scale structure.[39] A potentially disastrous difficulty with cryons is segregating them from luminous matter on small scales. Since this can only be done with non-gravitational forces (e.g., some form of dissipation), and since there is no evidence for dissipation on scales $\gtrsim$ galactic haloes this suggests that the cryon-to-luminous mass ratio should be constant on scales $\gtrsim$ galactic haloes--thereby implying that Ω is

$\ll 1$.

In fact in any scenario $\Omega = 1$ is problematic since there is no dynamic evidence that Ω is anywhere near as large as 1. Unless the dynamical evidence is misleading, the only way to reconcile it with $\Omega = 1$ is to uniformly spread out the mass density needed to bring Ω up to 1 (e.g. in relativistic particles or a relic cosmological constant) so that it would have escaped detection thus far.[40]

Let me end this section by reminding the reader of an important assumption which is almost always made when discussing structure in the Universe, that light is a good tracer of mass. This is a very strong assumption and if not true, could drastically alter some of the conclusions mentioned above.

THE EARLY UNIVERSE AS A HEP LABORATORY

In recent years the early Universe and various astrophysical environments have been used quite successfully to constrain the properties of particles which by virtue of their feeble interactions or large rest masses cannot be studied in terrestrial laboratories. Here I will go through 3 'case studies': neutrinos, axions, and superheavy magnetic monopoles.

Neutrinos - In the standard cosmology light neutrinos ($m \lesssim 1$ MeV) are always about as abundant as photons (today = 110 cm^{-3} per species), making them particularly easy to study. Any light neutrino species will make a significant contribution to the energy density of the Universe during the epoch of nucleosynthesis. Competition between the expansion rate (which is determined by the energy density) and the weak interaction rates determine the amount of ^{4}He synthesized. The ^{4}He abundance can thus be used to constrain the total number of light neutrino species[6]: $N_\nu \leq 4$. Stable neutrinos of any mass will contribute to the present mass density. The constraint on the mass density of the Universe, $\Omega h^2 \lesssim O(1)$, implies that any stable neutrino species be less massive than O(100 eV) or more massive than O(2 GeV). An unstable neutrino species which decays radiatively (e.g., $\nu \to \nu' + \gamma$ or $\nu \to \nu' + e^\pm$) runs the risk of 'photon pollution'. The photon spectrum of the Universe is shown in Fig. 7, and the mass-lifetime constraints which follow from insisting that neutrinos have negligible cosmological environmental impact are displayed in Fig. 8. For a mass-lifetime relationship of the form

$$\tau_\nu = 10^{-6} g (m_\mu/m_\nu)^5 \text{ sec}, \qquad (22)$$

where g is a numerical factor in the range 1-10^{12}, only neutrino masses less than O(100 eV) or more than O(10 MeV) are 'cosmologically/astrophysically safe'. [The quantity a is a model-dependent factor which takes into account mixing angles, GIM suppression factors, etc.] All of these neutrino constraints are reviewed in ref. 41 (including a complete list of references).

Axions - Models which employ the Peccei-Quinn[42] symmetry to solve the strong CP problem predict the existence of a light ($m_a \approx m_\pi^2/f_{PQ}$), pseudo Nambu-Goldstone boson[43] ('the axion') with coupling strength $\propto f_{PQ}^{-1}$ (f_{PQ} = scale of Peccei-Quinn SSB). Astrophysics and cosmology severely constrain the allowed values of f_{PQ}. If O(100 GeV) $\lesssim f_{PQ} \lesssim$ O(10^9 GeV) axions are so copiously produced in the cores of main sequence and red giant stars (where temperatures reach 1-10 keV) that they would carry off the bulk of the energy produced by nuclear reations in the star. Since the lifetime of a star is determined by how rapidly it can 'get rid' of its nuclear free energy, axion emission would hasten the evolution of these stars, shortening the lifetime so drastically that such stars would be too shortlived to observe.[44] [Laboratory experiments rule out $f_{PQ} \lesssim$ few 100 GeV.]

The troubles for the axion do not end here. Although a weakly-coupled ($f_{PQ} \gtrsim 10^9$ GeV) axion (dubbed the 'invisible axion'[45]) is astrophysically safe, it is not necessarily cosmologically safe, as was pointed out recently.[46] The potential which 'anchors' θ_a (the angular order parameter associated with the SSB of the PQ symmetry) at a CP-conserving value is due to instanton effects, and does not develop until a temperature of O(Λ_{QCD}). At higher temperatures there is nothing that picks out a particular value of θ_a, i.e., $V(\theta_a)$ is flat, and so one expects the initial misalignment of θ_a relative to the minimum of $V(\theta)$ to be of O(1). When the instanton effects turn on θ_a begins oscillating; the energy density of these coherent oscillations behaves like a very non-relativistic (cold) gas of axions. The energy density of these oscillations today corresponds to

$$\Omega_a h^2 \simeq 10^{-1}(f_{PQ}/10^{12}\ \mathrm{GeV})^{11/9}; \qquad (23)$$

therefore, unless f_{PQ} is $\lesssim 10^{13}$ GeV, axions will contribute too much mass density to the Universe. Astrophysics and cosmology leave but a small window for f_{PQ}: 10^9 GeV-10^{13} GeV. Of course, it goes without saying that the case of $f_{PQ} \simeq 10^{13}$ GeV is of great cosmological interest (see Table I).

Superheavy Magnetic Monopoles - The monopoles predicted to exist in GUTs like SU(5) have three very conspicuous properties: (1) macroscopic mass (10^{16} GeV =

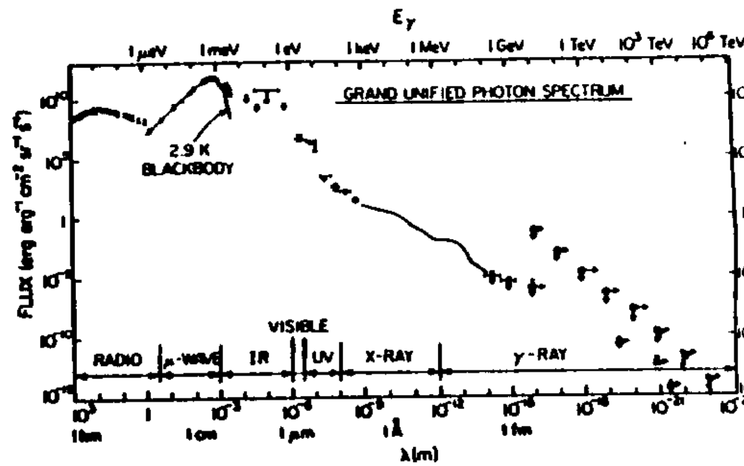


Figure 7 - The diffuse photon spectrum of the Universe from λ = 1 km to $\lambda = 10^{-24}$ m. Vertical arrows indicate upper limits; horizontal arrows indicate an integrated flux measurement ($E \gtrsim E_o$).

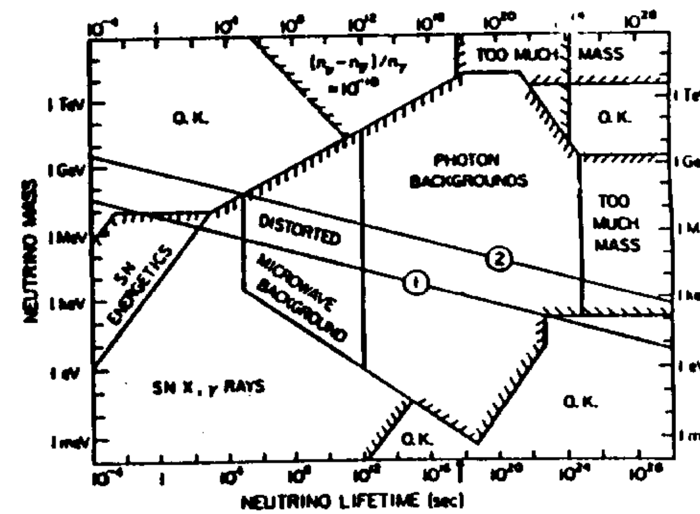


Figure 8 - Summary of astrophysical and cosmological constraints on neutrino masses and lifetimes. Lines marked 1 and 2 indicate mass-lifetime relationship for a = 1, 10^{12}, cf., eqn. (22).

10^{-8} gram in SU(5)); (2) hefty magnetic charge $g = (2e)^{-1} = 69\ e$; (3) the ability to catalyse nucleon decay with a strong interaction cross section:[47] $\sigma\beta = 10^{-28}\ cm^2$ (βc = nucleon-monopole relative velocity). These three properties lead to very stringent astrophysical/cosmological bounds on the relic flux of magnetic monopoles, which I will briefly review here (see Fig. 9).

The mass density in our neighborhood of the galaxy is $\approx 10^{-24}\ g\ cm^{-3}$--the monopole contribution cannot exceed this. This leads to an 'iron-clad' monopole flux limit of

$$F \lesssim 3 \times 10^{-10}(10^{16}\ GeV/m)\ cm^{-2}sr^{-1}s^{-1}. \qquad (24)$$

[Here and throughout I will assume that the monopole velocity is $\approx 10^{-3}$ c. Any initial velcity the monopole had at birth has long since been redshifted away. All contemporary acceleration mechanisms (galactic magnetic field, gravitational field of the galaxy, peculiar velocity relative to the Hubble flow, etc.) lead to a velocity $\approx 10^{-3}$c.] If monopoles do not cluster in the galaxy and instead are smoothly distributed in the cosmos, then their contribution to the mass density must be $\lesssim 2\rho_{crit} \approx 4 \times 10^{-29}\ g\ cm^{-3}$, resulting in the bound,

$$F \lesssim 10^{-14}(10^{16}\ GeV/m)\ cm^{-2}sr^{-1}s^{-1}. \qquad (25)$$

[The magnetic field of our galaxy will quickly eject all monopoles less massive than about 10^{19} GeV, suggesting that monopoles less massive than 10^{19} GeV should be smoothly distributed.[48]]

Magnetic monopoles will be accelerated by magnetic fields and thereby gain KE ($2 \times 10^4 eV\ cm^{-1}\ G^{-1}$). In the galaxy ($B \approx 3 \times 10^{-6}$ G; coherence length $\approx$ 300 pc $\approx 3.1 \times 10^{20}$) this amounts to $\approx 6 \times 10^{10}$ GeV per 300 pc. The 'NO FREE LUNCH PRINCIPLE' says that this energy gain must be compensated for by a loss in magnetic field energy. Thus a monopole flux F will lead to the decay of the galactic field in a time $\tau \approx (B^2/8\pi)/(gB4\pi F) \approx 10^8\ yrs\ (10^{-16}cm^{-2}sr^{-1}s^{-1}/F)$ [$4\pi FgB$ = rate (per volume) at which monopoles gain KE.] Arguing that the galactic dynamo can regenerate the magnetic field in a time no shorter than $O(10^8$ yrs) Parker obtained the bound[49]

$$F \lesssim 10^{-16}cm^{-2}sr^{-1}s^{-1}; \qquad (26)$$

for monopoles more massive than about 10^{17} GeV the gravitational effects (of the galaxy) must be taken into account and a slightly less stringent bound follows (see ref. 48). The 'Parker bound' can be evaded if the magnetic field of the galaxy is due to magnetic plasma

oscillations[48,50] (in which case the total energy oscillates between magnetic field energy and monopole KE). Although this scenario seems rather unlikely, it is a possibility, albeit exotic. In order for the oscillations to avoid Landau damping (in which case the Parker bound is valid), the monopole phase velocity must exceed the gravitational velocity dispersion in the galaxy ($\approx 10^{-3}$ c), which leads to a <u>lower bound</u> to the monopole flux,

$$F \gtrsim \tfrac{1}{4}\, mv^3(g\ell)^{-2},$$
$$\gtrsim 10^{-12}(m/10^{16}GeV)cm^{-2}sr^{-1}s^{-1}, \qquad (27)$$

where ℓ = 300 pc is the coherence length of the galactic magnetic field, and v is the gravitational velocity dispersion in the galaxy ($\approx 10^{-3}$ c). This exotic possibility predicts a flux which is beginning to be in conflict with ionization-type search experiments (see Fig. 9).

The 'Parker argument' has also been applied to the survival of intracluster magnetic fields, and results in a much more stringent limit.[51]

$$F \lesssim 10^{-18}cm^{-2}sr^{-1}s^{-1}. \qquad (28)$$

However, the existence of these fields is less well-established.

Callan and Rubakov[47] have shown that due to 's-wave sucking' monopoles should catalyse nucleon decay with a large cross section ($\approx 10^{-28}\ cm^2$), and the most stringent limit on the monopole flux follows from considering monopole catalysis of nucleon decay. The basic idea is very simple. Monopoles impinging upon various astrophysical objects (neutron stars[52], white dwarfs,[53] main sequence stars,[54] Jupiters,[55] the earth[55] <u>etc.</u>) will lose sufficient energy (through electronic interactions) to become captured. Once inside these objects, they catalyse nucleon decay, thereby releasing energy at a rate $\approx 10^{18}\ (\rho_{object}/10^{14}\ g\ cm^{-3})$ ergs s^{-1} <u>per monopole</u>; this energy is thermalized and then radiated from the surface of the object as soft x-rays (neutron stars), UV (white dwarfs), heat (Jupiter and the earth). The most stringent limit comes from neutron stars. For example, observations made by the Einstein x-ray Satellite constrain the number of monopoles inside the old radio pulsar PSR 1929+10 to be $\lesssim 10^{12}$; since the number of monopoles in PSR 1929+10 is proportional to the monopole flux, this can be translated into a flux bound of

$$F \lesssim 10^{-21}(c\beta/10^{-28}cm^2)^{-1}cm^{-2}sr^{-1}s^{-1}; \qquad (29)$$

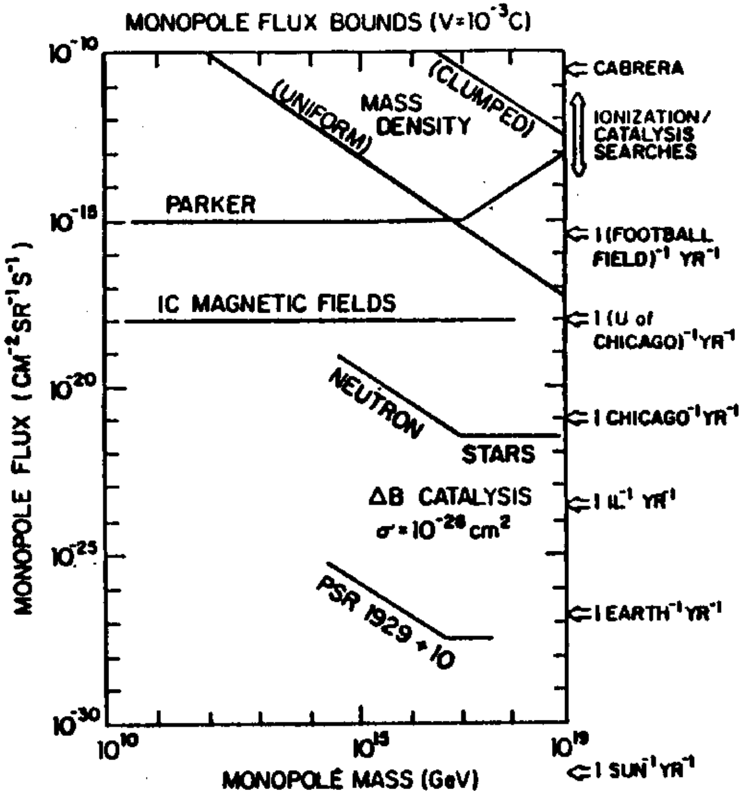


Figure 9 - Summary of the astrophysical and cosmological constraints on the monopole flux as a function of monopole mass. Where ever necessary the monopole velocity was taken to be $\simeq 10^{-3}$ c.

detection of such a tiny flux ($\simeq$ 1 Chicago^{-1} yr^{-1}) is a formidable (if not impossible) task! Flux bounds based upon measurements of the diffuse soft x-ray background and serendipitous searches for x-ray point sources give comparable limits. If the monopoles captured by the progenitor ($\simeq 10\ M_\odot$ star) of PSR 1929+10 are taken into account, the bound becomes:[54]

$$F \lesssim 2 \times 10^{-28} (\sigma\beta/10^{-28}\mathrm{cm}^2)^{-1} \mathrm{cm}^{-2}\mathrm{sr}^{-1}\mathrm{s}^{-1}, \qquad (30)$$

corresponding to slightly less than 1 earth^{-1} yr^{-1}. All of these monopole flux limits (summarized in Fig. 9) present a real challenge to the monopole hunter/huntress. These three case studies serve to illustrate how useful and powerful various astrophysical environments and the Universe itself are as non-standard HEP laboratories.

CONCLUDING REMARKS

Cosmology and particle physics both have very successful standard models: the hot big bang model which provides a reliable accounting of the history of the Universe from 10^{-2} sec until today, and the SU(3) x SU(2) x U(1) model which accurately describes particle physics at energies $\lesssim 10^3$ GeV. Progress in both fields has come to depend upon the link between the two, 'The Inner Space/Outer Space Connection'. The interplay between these two fields has already produced some very exciting and important results--baryogenesis, 'the monopole problem', the inflationary Universe paradigm, numerous candidates for the dark matter, and stringent constraints on neutrinos, axions, and monopoles. The future of research in this interdisciplinary field promises to be exciting, and I believe, will very likely provide us with some new surprises. Along this line let me conclude with some sobering thoughts about inflation. Inflation is extremely attractive because it offers the possibility of resolving a number of very fundamental cosmological puzzles with 'relatively well-known physics' (SSB transitions in gauge theories). However, any scenario which generates an enormous amount of entropy has the same potential (e.g., particle creation by quantum gravitational processes during the planck epoch, entropy creation due to compactification in a Kaluza-Klein theory, etc.). With regard to understanding the formation of structure in the Universe, strings[56] provide a potentially very attractive alternative to the creation of (adiabatic) density perturbations. The two key predictions of inflation, $\Omega = 1 \pm 10^{-\mathrm{BIG}\ \#}$ and the Zel'dovich spectrum of density perturbations, are so cosmologically-compelling, that I consider them to be a requirement for any viable candidate early Universe scenario (the ante to get into the game, so to speak), and if observationally verified cannot really be considered to be a verification of inflation.

This talk was meant to be a brief overview, rather than a thorough review, of an exciting and rapidly growing field (and was written under great duress!). I have tried to cite key papers and reviews which contain more complete lists of references, rather than to provide a thorough list of references. I apologize for my errors of omission, both in subject material and in the literature cited. This work was supported by the DOE both at Fermilab and Chicago (AC02-80ER-10773 A004), by NASA at Fermilab, and by the Alfred P. Sloan Foundation.

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