Fermi National Accelerator Laboratory

FERMILAB-Conf.-84/111-A
October, 1984

SUPERHEAVY MAGNETIC MONOPOLES AND THE STANDARD COSMOLOGY

Michael S. Turner

Theoretical Astrophysics Group
Fermi National Accelerator Laboratory
Batavia, Illinois 60510

and

Enrico Fermi Institute
The University of Chicago
Chicago, Illinois 60637

To appear in the Proceedings of Monopole '83 (ed, J. Stone; Plenum: NY, 1984)

SUPERHEAVY MAGNETIC MONOPOLES AND THE STANDARD COSMOLOGY

Michael S. Turner

Theoretical Astrophysics Group
Fermi National Accelerator Laboratory
Batavia, Illinois 60510

and

Enrico Fermi Institute
The University of Chicago
Chicago, Illinois 60637

ABSTRACT

The superheavy magnetic monopoles predicted to exist in grand unified theories (GUTs) are very interesting objects, both from the point of view of particle physics, as well as from astrophysics and cosmology. Astrophysical and cosmological considerations have proved to be invaluable in studying the properties of GUT monopoles. Because of the glut of monopoles predicted in the standard cosmology for the simplest GUTs (so many that the Universe should have reached a temperature of 3 K at the tender age of $\approx 10{,}000$ yrs), the simplest GUTs and the standard cosmology are not compatible. This is a very important piece of information about physics at unification energies ($E \gtrsim 10^{14}$ GeV) and about the earliest moments ($t \lesssim 10^{-34}$ s) of the Universe. In this talk I review the cosmological consequences of GUT monopoles within the context of the standard hot big bang model.

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} s 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.

Grand unification provides a particularly good example of the importance of the interplay between particle physics and cosmology. GUTs give us only a 'few windows' from our low energy world to the physics of unification energies, the two most familar and important ones being baryon number nonconservation and superheavy magnetic monopoles. Both of these predictions also have profound implications for cosmology and astrophysics. Baryon nonconservation provides for the first time a framework for understanding the 'baryon asymmetry of the Universe' (for a recent review see Ref. 1). Baryogenesis is a spectacular and unqualified success for the marriage of cosmology and GUTs, and a useful window to very high energies and the earliest history of the Universe. Monopoles, on the other hand, have been a cosmological disaster. In the standard cosmology so many monopoles should have been produced (at least for the simplest GUTs) that the Universe should have reached a temperature of 3 K while still in its infancy ($t \approx 10{,}000$ yrs). This is the so-called 'Monopole Problem': the simplest GUTs and the standard cosmology are not compatible (at least at times as early as 10^{-34} s). Although somewhat discouraging, this too is an important piece of information about physics at very high energies and the earliest history of the Universe. Monopoles have been a real boon for astrophysics. Because of their macroscopic masses [$\approx 10^{-8}$ g in SU(5)], hefty magnetic charge (integer multiple of ≈ 69 e), and their remarkable ability to catalyze nucleon decay at a prodigious rate, monopoles if they exist, should today be doing very interesting things--contributing mass density, 'shorting out' astrophysical magnetic fields, and gobbling up nucleons (releasing 1 GeV per gobble) to mention but three. Astrophysical considerations have resulted in very stringent bounds on the flux of relic monopoles. I have summarized these bounds in Fig. 1, and they will be discussed in great detail by other speakers.

In this talk I will first briefly review the standard cosmology. Next I will discuss monopole production in the very early Universe, both as topological defects (the 'Kibble mechanism') and by energetic particle collisions. The glut which results from the Kibble mechanism in the standard cosmology is the root of the monopole problem. On the other hand, if Kibble production can be suppressed, we are apparently left with a

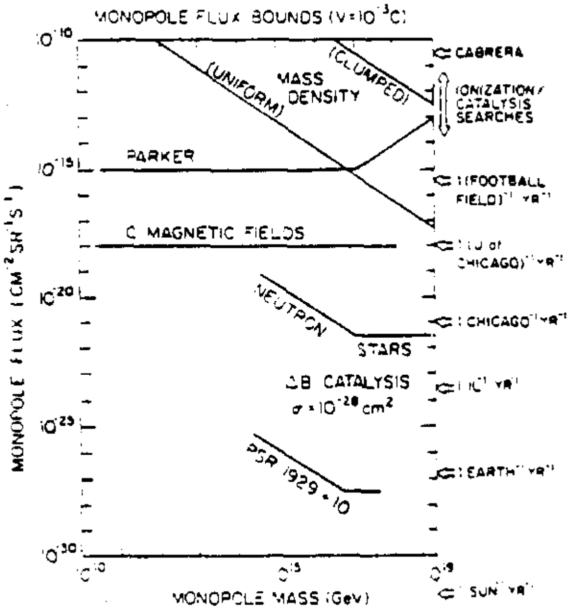

Fig. 1 Summary of the astrophysical and cosmological constraints on the monopole flux as a function of monopole mass. Wherever necessary the monopole velocity was taken to be $\approx 10^{-3}$ c. The magnetic field bounds are discussed in Refs. 2-5; the bounds based upon monopole catalysis of nucleon decay in neutron stars are discussed in Refs. 6-11. The most stringent bound (based upon monopole catalysis of nucleon decay in PSR 1929 + 10) was obtained also by taking into account the monopoles that the progenitor of this pulsar captured while it was on the main sequence.[8]

monopole famine, due to the dearth of monopoles produced in energetic particle collisions (although the uncertainties here are exponential). I will then trace the history of monopoles from their birth during the earliest moments of the Universe, through their adolescence, until today, with the aim of answering the important questions like, where should one expect to find monopoles today?, and with what velocities should they be moving? I will finish with some concluding remarks.

THE STANDARD COSMOLOGY[12]

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 ≈ 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 $t \approx 10^{-2}$ s - 10^2 s when the temperature was ≈ 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[13] (see Figs. 2 and 3):

$$\eta = (4 - 7) \times 10^{-10} \tag{1}$$

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

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

where the Hubble parameter $H_0 = 100\ h\ \mathrm{kms^{-1}\ Mpc^{-1}}$, and T is the present temperature of the cosmic microwave background. Observations strongly suggest that: $1/2 < h < 1$ and $2.7\ \mathrm{K} < T < 3.0\ \mathrm{K}$,

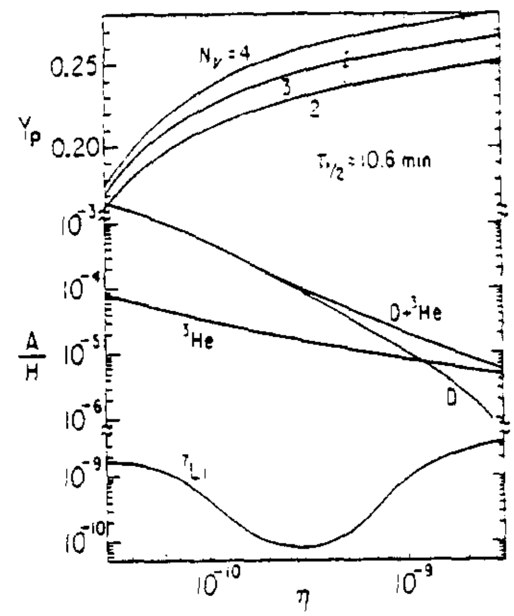

Fig. 2 The predicted primordial abundances of D, ^{3}He, ^{4}He, and ^{7}Li [$\tau_{1/2}(n) = 10.6$ min was used; error bar shows $\Delta\tau_{1/2} = \pm 0.2$ min; Y_p = mass fraction of ^{4}He]. Inferred primordial abundances: $Y_p \approx 0.23 - 0.25$; $(D/H)_p \gtrsim 10^{-5}$; $(D + {}^3He)_p/H \lesssim 10^{-4}$; $({}^7Li/H)_p \approx (1.1 \pm 0.4) \times 10^{-10}$. Consistency of the predicted abundances with observations can only be achieved for $\eta \approx (4 - 7) \times 10^{-10}$ ($\approx$ baryon-to-photon ratio) and $N_\nu \lesssim 4$ (= number of light neutrino species). For $4 < \eta/10^{-10} < 7$, $0.014 < \Omega_b < 0.15$. See Ref. 13 for more details.

so that the concordant range for η implies

$$0.014h^{-2}\ (T/2.7\ K)^3 < \Omega_b < 0.034h^{-2}\ (T/3.0\ K)^3, \tag{3a}$$

$$0.014 < \Omega_b < 0.15; \tag{3b}$$

implying that baryons alone cannot provide the closure density.

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 of ^{4}He ($\approx$ 25% by mass) and D (D/H $\approx$ few $\times 10^{-5}$) can be produced. It is the successful predictions of big bang nucleosynthesis that gives us confidence in the standard model back to $\approx 10^{-2}$ s after 'the bang'.

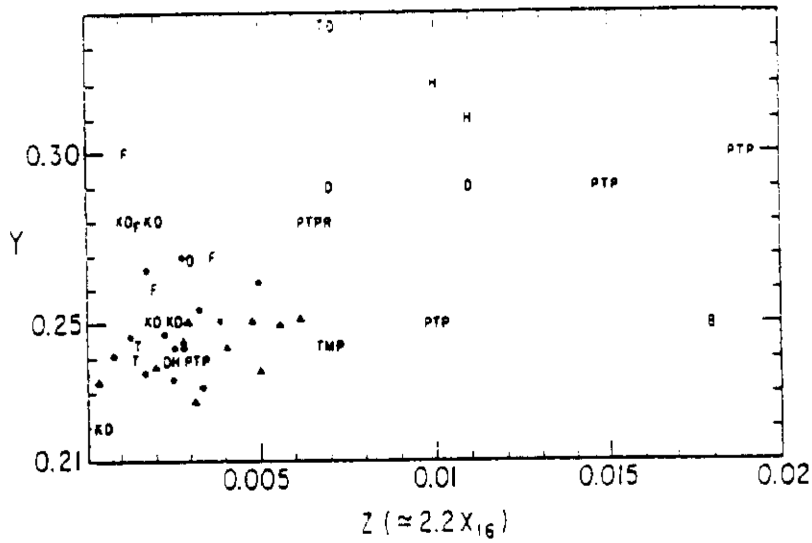

Fig. 3 Summary of determinations of ^{4}He mass fraction (Y) in HII regions as a function of the metal abundance Z (more precisely, 2.2 times the mass fraction of ^{16}O). Where the metal abundance is lowest, one expects the stellar contribution to Y to be the smallest. The data exhibit this trend and clearly show the existence of a primordial ^{4}He component of about 0.23-0.25 (by mass). For more details see Ref. 13.

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[12]

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

where ds^2 is the proper separation between two spacetime 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), or the wavelength of a freely-propagating photon ($\lambda \propto R$). 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 $100\ h\ \mathrm{kms}^{-1}\ \mathrm{Mpc}^{-1} = h\ (10^{10}\ \mathrm{yr})^{-1}$ $(1/2 < h < 1)$. As can be seen from Eqn. 5a model Universes with $k < 0$ 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:

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

$$\equiv \Omega - 1,$$

where $\Omega = \rho/\rho_{crit}$ and $\rho_{crit} = 3H^2/8\pi G \simeq 1.88\ h^2 \times 10^{-29}\ \mathrm{gcm}^{-3}$. From primordial nucleosynthesis we know that $\Omega \geq \Omega_b \gtrsim 0.014$. The best upper limit to Ω follows by considering the age of the Universe:

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

where $f(\Omega) < 1$ 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 (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 $\left(\text{recall } \lambda \propto R(t)\right)$. The energy density contributed by a relativistic species $(T >> m)$ at temperature T is

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

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 $h/2\pi = c = k_B = 1$, so that $1\ \mathrm{GeV} = (1.97 \times 10^{-14}\ \mathrm{cm})^{-1} = (1.16 \times 10^{13}\ \mathrm{K}) = (6.57 \times 10^{-25}\ \mathrm{s})^{-1}$, $G = m_{Pl}^{-2}$ $(m_{Pl} = 1.22 \times 10^{19}\ \mathrm{GeV})$, and $1\ \mathrm{GeV}^4 = 2.32 \times 10^{17}\ \mathrm{g\ cm}^{-3}$.

Today, the energy density contributed by relativistic particles (photons and 3 neutrino species) is negligible: $\Omega_{rel} = 4 \times 10^{-5}\ h^{-2}\ (T/2.7\ \mathrm{K})^4$. However, since $\rho_{rel} \propto R^{-4}$, while

$\rho_{nonrel} \propto R^{-3}$, early on relativistic species will dominate the energy density. For $R/R_{today} < 4 \times 10^{-5}\ (\Omega h^2)^{-1}\ (T/2.7\ K)^4$, which corresponds to $t < 4 \times 10^{10}$ s $(\Omega h^2)^{-2}\ (T/2.7\ K)^6$ and $T > 6$ eV $(\Omega h^2)(2.7\ K/T)^3$, 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:

$$H \equiv (\dot{R}/R) \simeq (4\pi^3\ g_*/45)^{1/2}\ T^2/m_{pl}, \qquad (9)$$
$$\simeq 1.66\ g_*^{1/2}\ T^2/m_{pl},$$

(valid for $t \lesssim 10^{10}$ s, $T \gtrsim 10$ eV).

Here g_* counts the total number of effective degrees of freedom of all the relativistic particles (i.e., those species with mass $\ll T$):

$$g_* = \sum_{Bose} g_i(T_i/T)^4 + 7/8 \sum_{Fermi} g_i(T_i/T)^4 \qquad (10)$$

and T is the photon temperature. For example: $g_*(3\ K) = 3.36$ (γ, $3\nu\bar{\nu}$); g_*(few MeV) = 10.75 (γ, $e^\pm$, 3 $\nu\nu$); g_*(few 100 GeV) = 110 (γ, $W^\pm$ $Z°$, 8 gluons, 3 families of quarks and leptons, and 1 Higgs doublet).

If thermal equilibrium is maintained, then the second Friedmann equation, Eqn. 5b - conservation of energy, implies that the entropy per comoving volume (a volume with fixed r, θ, ϕ coordinates) $S \propto sR^3$ remains constant. Here s is the entropy density, which is dominated by the contribution from relativistic particles, and

$$s = (\rho + p)/T = 2\pi^2\ g_*\ T^3/45, \qquad (11)$$

which is proportional to the number density of relativistic particles. So long as the expansion is adiabatic (i.e., in the absence of entropy production) S(and s) will prove to be useful fiducials. For example, at low energies ($E \ll 10^{14}$ GeV) baryon number is effectively conserved, and so the net baryon number per comoving volume $N_B \propto n_B(\equiv n_b - n_{\bar{b}})\ R^3$ remains constant, implying that the ratio n_B/s is a constant of the expansion. Today $s \simeq 7n_\gamma$, so that

$$n_B/s \simeq \eta/7 \simeq (6 - 10) \times 10^{-11}. \qquad (12)$$

Once monopole-antimonopole annihilations are no longer important, the number of monopoles per comoving volume is also conserved, $N_M \propto n_M\ R^3$. Comparing this to the baryon number and entropy per comoving volume we get two ratios which remain constant (so long as annihilations are not important, and entropy and baryon number are conserved) and are related to the present average flux of monopoles in the Universe by:

$$\langle F_M \rangle = 10^{10}\ (n_M/s)(v_M/10^{-3}\ c)\mathrm{cm}^{-2}\ \mathrm{sr}^{-1}\ \mathrm{s}^{-1}, \tag{13a}$$

$$= (n_M/n_B)(v_M/10^{-3}\ c)\mathrm{cm}^{-2}\ \mathrm{sr}^{-1}\ \mathrm{s}^{-1}, \tag{13b}$$

where v_M is the typical monopole velocity (which will be discussed at length later). The fraction of critical density contributed by monopoles (Ω_M) is:

$$\Omega_M = 10^{24}\ (n_M/s)(m_M/10^{16}\ \mathrm{GeV}), \tag{14a}$$

$$= 10^{14}\ \langle F_M \rangle (10^{-3}\ c/v_M)\mathrm{cm}^2\ \mathrm{sr\ s}. \tag{14b}$$

Whenever $g_* \simeq$ constant, the constancy of the entropy per comoving volume implies that $T \propto R^{-1}$; together with Eq. 9 this gives

$$R(t) = R(t_o)(t/t_o)^{1/2}, \tag{15a}$$

$$t \simeq 0.3\ g_*^{-1/2}\ m_{pl}/T^2,$$

$$\simeq 2.4 \times 10^{-6}\ \mathrm{s}\ g_*^{-1/2}\ (T/\mathrm{GeV})^{-2}, \tag{15b}$$

valid for $t \lesssim 10^{10}$ s and $T \gtrsim 10$ eV.

Finally, let me mention one more important feature of the standard cosmology, the existence of particle horizons. The distance that a light signal could have propagated since the bang is finite, and easy to compute. Photons travel on paths characterized by $ds^2 = 0$; for simplicity (and without loss of generality) consider a trajectory with $d\theta = d\phi = 0$. The coordinate distance covered by this photon since 'the bang' is just $\int_0^t dt'/R(t')$, corresponding to a physical distance (measured at time t) of

$$d_H(t) = R(t) \int_0^t dt'/R(t') \tag{16}$$

$$= t/(1 - n)\ [\text{for } R \propto t^n,\ n < 1]$$

If $R \propto t^n$, then this distance is finite and $\simeq t \simeq H^{-1}$. Note that even if $d_H(t)$ diverges and there is no finite particle horizon (e.g., if $R \propto t^n$, $n > 1$), the Hubble radius H^{-1} still sets the scale for the 'physics horizon'. Since all physical length scales roughly e-fold in a time H^{-1}, causally-coherent microphysical processes can only operate over times (also distances) $\lesssim H^{-1}$.

During the radiation-dominated epoch $n = 1/2$ and $d_H = 2t$; the baryon number and entropy within the horizon at time t are easily computed:

$$S_{HOR} = (4\pi/3)t^3\ s,$$

$$\simeq 0.05\ g_*^{-1/2}\ (m_{pl}/T)^3; \tag{17a}$$

$$N_{B-HOR} = (n_B/s) \times S_{HOR},$$

$$\simeq 10^{-12}\ (m_{pl}/T)^3; \tag{17b}$$

note that I have implicitly assumed the constancy of the baryon-to-entropy ratio in computing N_{B-HOR}.

Although our verifiable knowledge of the early history of the Universe only takes us back to $t \simeq 10^{-2}$ s, and $T \simeq 10$ MeV, nothing in our present understanding of the laws of physics suggests that it is unreasonable to extrapolate back to times as early as $\simeq 10^{-43}$ s and temperatures as high as $\simeq 10^{19}$ 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 Fig. 4.

BIRTH: GLUT OR FAMINE

In 1931 Dirac[14] showed that <u>if</u> magnetic monopoles exist, then the single-valuedness of quantum mechanical wavefunctions requires the magnetic charge of a monopole to satisfy the quantization condition

$$g = ng_D,\ n = 0,\ \pm1,\ \pm2\ \ldots$$

$$g_D = 1/2e \simeq 69e.$$

However, one is <u>not</u> required to have Dirac monopoles in the theory--you can take 'em or leave 'em! In 1974 't Hooft[15] and Polyakov[16] independently made a remarkable discovery. They showed that monopoles are <u>obligatory</u> in the low-energy theory whenever a semi-simple group G, e.g., SU(5); breaks down to a group G' × U(1) which contains a U(1) factor [e.g., SU(3) × SU(2) × 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 $m_M \simeq O(M/\alpha)$ [$\simeq 10^{16}$ GeV in SU(5); M = scale of SSB], and have a magnetic

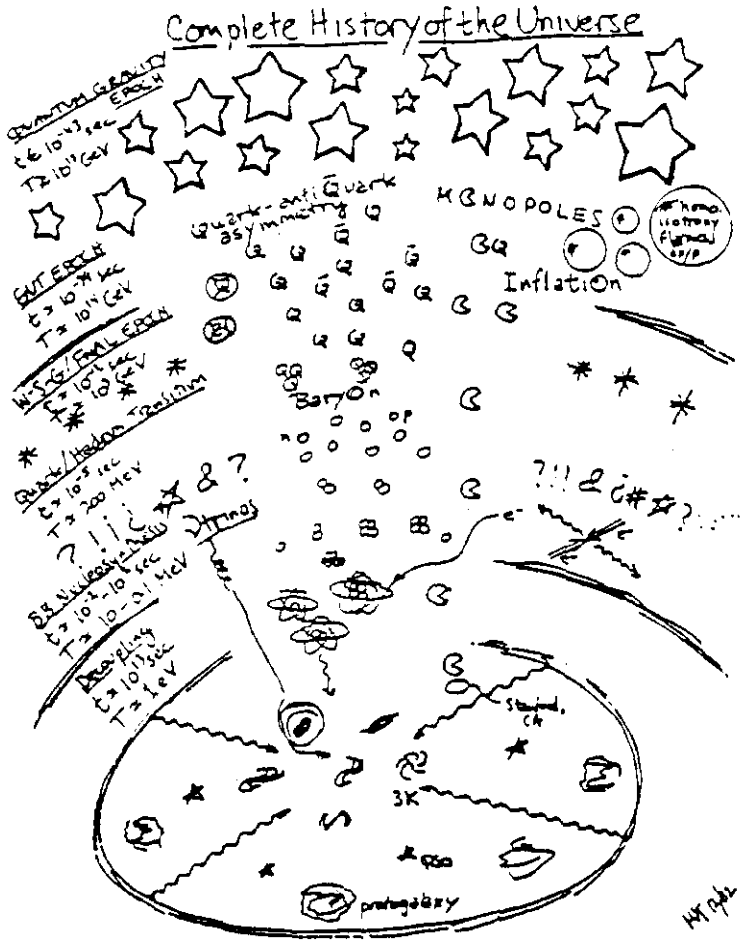

Fig. 4 'The *Complete History of the Universe*'. Highlights include: decoupling ($t \approx 10^{13}$ s, $T \approx 1/3$ eV) - the surface of last scattering for the cosmic microwave background, epoch after which matter and radiation cease to interact and matter 'recombines' into neutral atoms (D, ^{3}He,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

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} s after 'the bang' when the temperature was $\simeq O(10^{14}$ 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[17], Preskill[18], and Zel'dovich and Khlopov[19], and I will 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):

	$\uparrow$	
$\uparrow\ \uparrow\ \uparrow$	$\leftarrow\quad\rightarrow$	$\rightarrow$ = direction of
$\uparrow\ \uparrow\ \uparrow$	$\downarrow$	Higgs field
<u>no monopole</u>	<u>monopole</u>	in group space.

Clearly monopole configurations cannot exist until the SSB $[G \rightarrow G' \times U(1)]$ transition takes place. When spontaneous symmetry breaking occurs, the Higgs field can only be smoothly oriented (i.e., the no monopole configuration) on scales smaller than some characteristic correlation length ξ. On the microphysical side, the inverse Higgs mass at the Ginzburg temperature (T_G) sets such a scale: $\xi \simeq m_H^{-1}(T_G)$ (in a second-order phase transition)[20]. [The Ginzburg temperature is the temperature at which it becomes improbable for the Higgs field to fluctuate between the SSB minimum and $\phi = 0$.] Cosmological considerations set an absolute upper bound: $\xi \lesssim d_H(\simeq t$ in the standard cosmology). [Note, even if the $d_H(t)$ diverged, e.g., because $R \propto t^n$ $(n > 1)$ for $t \lesssim t_{pl}$, the physics horizon H^{-1} sets an absolute upper bound on ξ, which is numerically identical.] On scales larger than ξ the Higgs field must be uncorrelated, and thus we expect of order 1 monopole per correlation volume ($\simeq \xi^3$) to be produced as a topological defect when the Higgs field freezes out.

Let's focus on the case where the phase transition is either second order or weakly-first order. Denote the critical temperature for the transition by T_c $\left(\simeq O\ (M)\right)$, and as before the monopole mass by $m_M \simeq O(M/\alpha)$. The age of the Universe when $T = T_c$

is given in the standard cosmology by: $t_c \simeq 0.3\, g_*^{-1/2}\, m_{pl}/T_c^2$, cf. Eqn. 15b. For SU(5): $T_c \simeq 10^{14}$ GeV, $m_M \simeq 10^{16}$ GeV and $t_c \simeq 10^{-34}$ s. Due to the fact that freezing of the Higgs must be uncorrelated on scales $\gtrsim \xi$, we expect an initial monopole abundance of O(1) per correlation volume; using $d_H(t_c)$ as an absolute upper bound on ξ this leads to: $(n_M)_i \simeq O(1)\, t_c^{-3}$. Comparing this to our fiducials S_{HOR} and N_{B-HOR}, we find that the initial monopole-to-entropy and monopole-to-baryon number ratios are:

$$n_M/s \gtrsim 10^2\, (T_c/m_{pl})^3, \tag{18a}$$

$$n_M/n_B \gtrsim 10^{12}\, (T_c/m_{pl})^3. \tag{18b}$$

Preskill[18] has shown that unless n_M/s is $> 10^{-10}$ monopole-antimonopole annihilations do not significantly reduce the initial monopole abundance. If $n_M/s > 10^{-10}$, he finds that n_M/s is reduced to $\simeq 10^{-10}$ by annihilations. For $T_c < 10^{15}$ GeV our estimate for n_M/s is $< 10^{-10}$, and we will find that in the standard cosmology T_c must be $\ll 10^{15}$ GeV to have an acceptable monopole abundance, so for our purposes we can ignore annihilations. Assuming that the expansion has been adiabatic since $T \simeq T_c$, this estimate for n_M/s translates into:

$$\langle F_M\rangle \simeq 10^{-3}\, (T_c/10^{14}\ \mathrm{GeV})^3\ \mathrm{cm}^{-2}\ \mathrm{sr}^{-1}\ \mathrm{s}^{-1}, \tag{19a}$$

$$\langle \Omega_M\rangle \simeq 10^{11}\, (T_c/10^{14}\ \mathrm{GeV})^3 (m_M/10^{16}\ \mathrm{GeV}) \tag{19b}$$

--a flux that would make any monopole hunter/huntress ecstatic, and an Ω_M that is unacceptably large (except for $T_c \ll 10^{14}$ GeV). As was discussed previously, Ω can be at most O(few), so we have a very big problem with the simplest GUTs (in which $T_c \simeq 10^{14}$ GeV). This is the so-called 'Monopole Problem'. The statement that $\Omega_M \simeq 10^{11}$ for $T_c \simeq 10^{14}$ GeV is a bit imprecise; clearly if $k < 0$ (corresponding to $\Omega < 1$) monopole production cannot close the Universe (and in the process change the geometry from being infinite in extent and negatively-curved, to being finite in extent and positively-curved). More precisely, a large monopole abundance would result in the Universe becoming matter-dominated much earlier, at $T \simeq 10^3$ GeV $(T_c/10^{14}\ \mathrm{GeV})^3$ $(m_M/10^{16}\ \mathrm{GeV})$, and eventually reaching a temperature of 3 K at the young age of $t \simeq 10^4$ yrs $(T_c/10^{14}\ \mathrm{GeV})^{-3/2}$ $(m_M/10^{16}\ \mathrm{GeV}^{-1/2}$. The requirement that $\Omega_M \lesssim$ O(few) implies that

$$T_c \lesssim 10^{11}\ \mathrm{GeV} \qquad (\Omega_M \lesssim \mathrm{few})$$

where I have taken m_M to be O(100 T_c). Note, given our generous estimate for ξ, even this is probably not safe; if one had a GUT in which $T_c \simeq 10^{11}$ GeV a more careful estimate for ξ would be called for.

The Parker bound[2-4] (see Fig. 1) on the average monopole flux in the galaxy, $\langle F_M \rangle \lesssim 10^{-15}\ \mathrm{cm}^{-2}\ \mathrm{sr}^{-1}\ \mathrm{s}^{-1}$, results in a slightly more stringent constraint:

$$T_c \lesssim 10^{10}\ \mathrm{GeV} \qquad \text{(Parker bound)}$$

The most restrictive constraints on T_c follow from the neutron star catalysis bounds on the monopole flux[6-11], and the most restrictive of those[8], $\langle F_M \rangle \lesssim 10^{-27}\ \mathrm{cm}^{-2}\ \mathrm{sr}^{-1}\ \mathrm{s}^{-1}$, implies that

$$T_c \lesssim 10^{5}\ \mathrm{GeV} \qquad \text{(Neutron star catalysis bound)}$$

Note, to obtain these bounds I have compared my estimate for the average monopole flux in the Universe, Eqn. 19a, with the astrophysical bounds on the average flux of monopoles in our galaxy. If monopoles cluster in galaxies (which I will later argue is unlikely), then the average galactic flux of monopoles is greater than the average flux of monopoles in the Universe, making the above bounds on T_c <u>more</u> restrictive.

If the GUT transition is strongly first order (I am excluding inflationary Universe scenarios for the moment), then the transition will proceed by bubble nucleation at a temperature T_n ($\ll T_c$), when the nucleation rate becomes comparable to the expansion rate H. Within each bubble the Higgs field is correlated; however, the Higgs field in different bubbles should be uncorrelated. Thus one would expect O(1) monopole per bubble to be produced. When the Universe supercools to a temperature T_n, bubbles nucleate, expand, and rapidly fill all of space; if r_b is the typical size of a bubble when this occurs, then one expects n_M to be $\simeq r_b^{-3}$. After the bubbles coalesce, and the Universe reheats, the entropy density is once again $s \simeq g_* T_c^3$, so that the resulting monopole to entropy ratio is: $n_M/s \simeq (g_* r_b^3 T_c^3)^{-1}$. Guth and Weinberg[21] have calculated r_b and find that $r_b \simeq (m_{pl}/T_c^2)/\ln(m_{pl}^4/T_c^4)$, leading to a relatively accurate estimate for the monopole abundance:

$$n_M/s \simeq [\ln(m_{pl}^4/T_c^4)(T_c/m_{pl})]^3,$$

which is even more disasterous than the estimate for a second order phase transition [recall, however, estimate (18) was an absolute <u>lower</u> bound].

The bottom line is that we have a serious problem here--the standard cosmology extrapolated back to $T \simeq T_c$ and the simplest GUTs are incompatible (to say the least). One (or both) must be modified. This is a valuable piece of information.

A number of possible solutions have been suggested. To date the most attractive is the new inflationary Universe scenario (which will be discussed at length by P. Steinhardt). In this scenario, a small region (size $\lesssim$ the horizon) within which the Higgs field could be correlated, grows to a size which encompasses all of the presently observed Universe, due to the exponential expansion which occurs during the phase transition. This results in less than one monopole in the entire observable Universe (due to Kibble production).

Let me very briefly review some of the other attempts to solve the monopole problem. Several people have pointed out that if there is no complete unification [e.g., if $G = H \times U(1)$], or if the full symmetry of the GUT is not restored in the very early Universe (e.g., if the maximum temperature the Universe reached was $< T_c$, or if a large lepton number[22], $n_L/n_\gamma > 1$, prevented symmetry restoration at high temperatures), then there would be no monopole problem. However, none of these possibilities seems particularly attractive.

Several authors[23-26] have studied the possibility that monopole-antimonopole annihilation could be enhanced over Preskill's estimate, due to 3-body annihilations or the gravitational clumping of monopoles (or both). Thus far, this approach too has been unsuccessful.

Bais and Rudaz[27] have suggested that large fluctuations in the Higgs field at temperatures near T_c could allow the monopole density to relax to an acceptably small value. They do not explain how this mechanism can produce the acausal correlations needed to do this.

Mechanisms have been proposed in which monopoles and antimonopoles form bound pairs connected by flux tubes, leading to rapid monopole-antimonopole annihilation. For example, Linde[28] proposed that at high temperatures color magnetic charge is confined, and Lazarides and Shafi[29] proposed that monopoles and antimonopoles become connected by Z° flux tubes after the $SU(2) \times U(1)$ SSB phase transition. In both cases, however, the proposed flux tubes are not topologically stable, nor has their existence even been demonstrated.

Langacker and Pi[30] have proposed a solution which does seem to work. It is based upon an unusual (and perhaps contrived) symmetry breaking pattern for SU(5):

$$SU(5) \underset{T_c \simeq 10^{14}\ \mathrm{GeV}}{\rightarrow} SU(3) \times SU(2) \times U(1) \underset{T_1}{\rightarrow} SU(3) \underset{T_2}{\rightarrow} SU(3) \times U(1)$$

$$\underbrace{T_1 \longleftrightarrow T_2}_{\text{superconducting phase}}$$

(note T_1 could be equal to T_c). The key feature of their scenario is the existence of the epoch ($T = T_1 \to T_2$) in which the U(1) of electromagnetism is sponteneously broken (a superconducting phase); during this epoch magnetic flux must be confined to flux tubes, leading to the annihilation of the monopoles and antimonopoles which were produced earlier on, at the GUT transition. Although somewhat contrived, their scenario appears to be viable (however, I'll have more to say about it shortly).

Finally, one could invoke the Tooth Fairy (in the guise of a perfect annihilation scheme). E. Weinberg[31] has recently made a very interesting point regarding 'perfect annihilation schemes', which applies to the Langacker-Pi scenario[30], and even to a Tooth Fairy which operates causally. Although the Kibble mechanism results in equal numbers of monopoles and antimonopoles being produced, E. Weinberg points out that in a finite volume there can be magnetic charge fluctuations. He shows that if the Higgs field 'freezes out' at $T = T_c$ and is uncorrelated on scales larger than the horizon at that time, then the expected net RMS magnetic charge in a volume V which is much bigger than the horizon is

$$\Delta n_M \simeq (V/t_c{}^3)^{1/3}. \tag{20}$$

He then considers a perfect, causal annihilation mechanism which operates from $T = T_1 \to T_2$ (e.g., formation of flux tubes between monopoles and antimonopoles). At best, this mechanism could reduce the monopole abundance down to the net RMS magnetic charge contained in the horizon at $T = T_2$, leaving a final monopole abundance of

$$n_M/s \simeq 10^2\ T_c\ T_2{}^2/m_{pl}{}^3, \tag{21}$$

resulting in

$$\Omega_M \gtrsim 0.1\ (T_c/10^{14}\ \mathrm{GeV})(m_M/10^{16}\ \mathrm{GeV})(T_2/10^8\ \mathrm{GeV})^2, \tag{22a}$$

$$\langle F_M \rangle \gtrsim 10^{-15}\ (T_c/10^{14}\ \mathrm{GeV})(T_2/10^8\ \mathrm{GeV})^2\ \mathrm{cm}^{-2}\ \mathrm{sr}^{-1}\ \mathrm{s}^{-1}. \tag{22b}$$

It is difficult to imagine a perfect annihilation mechanism which could operate at temperatures $\lesssim 10^3$ GeV, without having to modify the standard SU(2) × U(1) electroweak theory; for $T_c \simeq 10^{14}$GeV and $T_2 \simeq 10^3$ GeV, E. Weinberg's argument[31] implies that $\langle F_M \rangle$ must be $\gtrsim 10^{-25}\ \mathrm{cm}^{-2}\ \mathrm{sr}^{-1}\ \mathrm{s}^{-1}$, which would be in conflict with the most stringent neutron star catalysis bound[8], $F_M < 10^{-27}\ \mathrm{cm}^{-2}\mathrm{sr}^{-1}\ \mathrm{s}^{-1}$.

In a recent preprint A. Vilenkin[32] disputes E. Weinberg's argument about the significance of magnetic charge fluctuations, and provides several counterexamples. In an even more recent preprint Lee and E. Weinberg[33] refute Vilenkin's counterexamples, and provide the results of numerical simulations which support E.

Weinberg's original conclusions[31]. Clearly, this important issue is not settled yet. It does, however, seem clear that a perfect annihilation mechanism which operates down to a temperature T_2 can do no better than to reduce the monopole abundance to 1 per horizon volume at $T = T_2$, or $n_M/s \simeq 10^2\ (T_2/m_{pl})^3$--to do better would require Higgs field correlations on scales larger than the horizon. Therefore, at the very best, the T_c's in Eqns. 22a,b could be replaced by T_2's.

Finally, I should emphasize that the estimate of n_M/s based upon $\xi \lesssim d_H(t)$ is an absolute (and very generous) lower bound to n_M/s. Should a model be found which succeeds in suppressing the monopole abundance to an acceptable level (e.g., by having $T_c \ll 10^{14}$ GeV or by a perfect annihilation epoch), then the estimate for ξ must be refined and scrutinized.

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) + antiparticles(s) $\rightarrow$ 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 $\ll O(10^{10}$ GeV).] The numbers produced are intrinsically small because monopole configurations do not exist in the theory until SSB occurs ($T_c \simeq M \simeq$ scale of SSB), and have a mass $O(M/\alpha) \simeq 100\ M \simeq 100\ T_c$. For this reason they are never present in equilibrium numbers; however, some are produced due to the rare collisions of particles with sufficient energy. This results in a present monopole abundance of[34-36]

$$n_M/s \simeq 10^2\ (m_M/T_{max})^3 \exp(-2m_M/T_{max}), \tag{23a}$$

$$\Omega_M \simeq 10^{26}\ (m_M/10^{16}\ \text{GeV})(m_M/T_{max})^3 \exp(-2m_M/T_{max}), \tag{23b}$$

$$\langle F_M \rangle \simeq 10^{12}\ \text{cm}^{-2}\ \text{sr}^{-1}\ \text{s}^{-1}(m_M/T_{max})^3 \exp(-2m_M/T_{max}), \tag{23c}$$

where T_{max} is the highest temperature reached after SSB.

In general, $m_M/T_{max} \simeq O(100)$ so that $\Omega_M \simeq O(10^{-40})$ and $\langle F_M \rangle \simeq O(10^{-32}\ \text{cm}^{-2}\ \text{sr}^{-1}\ \text{s}^{-1})$--a negligible number of monopoles. However, the number produced is exponentially sensitive to m_M/T_{max}, so that a factor of 3-5 uncertainty in m_M/T_{max} introduces an enormous uncertainty in the predicted production. For example, in the new inflationary Universe, the monopole mass can be $\propto$ the Higgs field responsible for SSB, and as that field oscillates about the SSB minimum during the reheating process m_M also oscillates, leading to enhanced monopole production [m_M/T_{max} in Eqns. 23a, b, c is replaced by fm_M/T_{max}, where $f < 1$ depends upon the details of reheating; see Refs. 37, 38].

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!

FROM BIRTH THROUGH ADOLESCENCE ($t \simeq 10^{-34}$ s to $t \simeq 3 \times 10^{17}$ s)

As mentioned in the previous section, monopoles and antimonopoles do not annihilate in significant numbers; however, they do interact with the ambient charged particles (e.g., monopole + $e^- \rightleftarrows$ monopole + e^-) and thereby stay in kinetic equilibrium ($KE \simeq 3T/2$) until the epoch of $e^\pm$ annihilations ($T \simeq 1/2$ MeV, $t \simeq 10$ s). At the time of $e^\pm$ annihilations monopoles and antimonopoles should have internal velocity dispersions of:

$$\langle v_M^2\rangle^{1/2} \simeq 30 \text{ cm s}^{-1} (10^{16} \text{ GeV}/m_M)^{1/2}.$$

After this monopoles are effectively collisionless, and their velocity dispersion decays $\propto R(t)^{-1}$, so that if we neglect gravitational and magnetic effects, today they should have an internal velocity dispersion of

$$\langle v_M^2\rangle^{1/2} \simeq 10^{-8} \text{ cm s}^{-1} (10^{16} \text{ GeV}/m_M)^{1/2}.$$

Since they are collisionless, only their velocity dispersion can support them against gravitational collapse. With such a small velocity dispersion to support them they are gravitationally unstable on all scales of astrophysical interest ($\lambda_{Jeans} \simeq 10^{-8}$ LY).

After decoupling ($T \simeq 1/3$ eV, $t \simeq 10^{13}$ s) [or the epoch of matter domination in scenarios where the mass of the Universe is dominated by a nonbaryonic component], matter can begin to clump, and structure can start to form. Monopoles, too, should clump and participate in the formation of structure. However, since they cannot dissipate their gravitational energy, they cannot collapse into the more condensed objects (such as stars, planets, the disk of the galaxy, etc.) whose formation clearly must have involved the dissipation of gravitational energy. Thus, one would only expect to find monopoles in structures whose formation did not require dissipation (such as clusters of galaxies, and galactic haloes). However, galactic haloes are not likely to be a safe haven for monopoles in galaxies with magnetic fields; monopoles less massive than about 10^{20} GeV will, in less than 10^{10} yrs, gain sufficient KE from a magnetic field of strength a few $\times 10^{-6}$ G to reach escape velocity[4]. So we are led to the conclusion that initially monopoles should either be uniformly distributed through the cosmos, or clumped in clusters of galaxies or in the haloes of galaxies with weak or non-existent magnetic fields. Since our own galaxy has a magnetic field of strength $\approx$ few $\times 10^{-6}$ G, and is not a member of a cluster of galaxies, we would expect the local flux of monopoles to be not too different from the average monopole flux in the Universe.

Although monopoles initially have a very small internal velocity dispersion, there are many mechanisms for increasing their velocities. First, typical peculiar velocities (i.e., velocities relative to the Hubble flux) are $O(10^{-3}$ c), leading to a typical monopole-galaxy velocity of 10^{-3} c. Monopoles will be accelerated by the gravitational fields of galaxies (to $\approx 10^{-3}$ c $\approx$ orbital velocity in the galaxy), and if they encounter them, clusters of galaxies (to $\approx 3 \times 10^{-3}$ c). A typical monopole, however, will never encounter a galaxy or a cluster of galaxies, the respective mean free paths being: $L_{gal} \approx 10^{26}$ cm ($\approx 10^{-2}$ c $\times$ age of the Universe) and $L_{cluster} \approx 3 \times 10^{28}$ cm.

Monopoles will also be accelerated by magnetic fields. The intragalactic magnetic field strength is $\lesssim 3 \times 10^{-11}$ G (Ref.39), and results in a monopole velocity of

$$v_M = 3 \times 10^{-4}\ c\ (B/10^{-11}\ G)(10^{16}\ GeV/m_M).$$

The galactic magnetic field will accelerate monopoles in our galaxy to velocities of[4]

$$v_M = 3 \times 10^{-3}\ c\ (10^{16}\ GeV/m_M)^{1/2}.$$

Taking all of these 'sources of velocity' into account, we can make an educated estimate of the typical monopole-detector relative velocity.

From Table 1 below it should be clear that the typical monopole should be moving with a velocity of <u>at least</u> a few x 10^{-3} c with respect to an earth-based detector. It goes without saying that 'this fact' is an important consideration for detector design.

Although planets, stars, etc. should be monopole-free at the time of their formation, they will accumulate monopoles during their lifetimes. The number captured by an object is where M, R

$$N_M = (4\pi R^2)(\pi - sr)(1 + 2GM/Rv_M{}^2)\langle F_M\rangle \epsilon\tau, \tag{24}$$

and τ are the mass, radius and age of the object, v_M is the monopole velocity, and ϵ is the efficiency with which the object stops monopoles which strike its surface. The efficiency of capture ϵ depends upon the mass and velocity of the monopole, and

Table 1. Typical Monopole-Detector Relative Velocities

DETECTOR VELOCITY		MONOPOLE VELOCITY	
orbit in galaxy	$2/3 \times 10^{-3}$ c	galactic B-field	3×10^{-3} c $(10^{16}\ \mathrm{GeV}/m_M)^{1/2}$
orbit in solar system	10^{-4} c	grav. acceleration by galaxy	10^{-3} c
		grav. acceleration by sun	10^{-4} c
		monopole-galaxy relative velocity	10^{-3} c

its rate of energy loss in the object. The quantity $(1 + 2GM/R\, v_M^2)$ is just the ratio of the capture cross section to the geometric cross section. Main sequence stars of mass $(0.6 - 30)M_\odot$ will capture monopoles less massive than about 10^{18} GeV with velocities $\lesssim 10^{-3}$ c with good efficiency ($\varepsilon \simeq 1$); in its main sequence lifetime a star will capture approximately 10^{24} F_{-16} monopoles[40] (essentially independent of its mass). Here $\langle F_M \rangle = F_{-16}\ 10^{-16}\ \mathrm{cm}^{-2}\ \mathrm{sr}^{-1}\ \mathrm{s}^{-1}$. Neutron stars will capture monopoles less massive than about 10^{20} GeV with velocities $\lesssim 10^{-3}$ c with unit efficiency, capturing about 10^{21} F_{-16} monopoles in 10^{10} yrs. Planets like Jupiter can stop monopoles less massive than about 10^{16} GeV with velocities $\lesssim 10^{-3}$ c, accumulating about 10^{22} F_{-16} monopoles in 10^{10} yrs.[41]. A planet like the earth can only stop light or slowly-moving monopoles (for $m_M = 10^{16}$ GeV, v_M must be $\lesssim 3 \times 10^{-5}$ c). Once inside, monopoles can do interesting things, like catalyze nucleon decay, which keeps the object hot (leading to a potentially observable photon flux), and eventually depletes the object of all its nucleons. A monopole flux of $F_{-21}\ 10^{-21}\ \mathrm{cm}^{-2}\ \mathrm{sr}^{-1}\ \mathrm{s}^{-1}$ will cause a neutron star to evaporate in 10^{11} $F_{-21}^{-1/2}$ yrs, a Jupiter-like planet to evaporate in 5×10^{15} $F_{-21}^{-1/2}$ yrs, and an Earth-like planet to evaporate in 10^{18} $F_{-21}^{-1/2}$ yrs[42]. Accretion of monopoles by astrophysical objects, however, does not significantly reduce the monopole flux; the mean free path of a monopole in the galaxy is $\simeq 10^{42}$ cm.

CONCLUDING REMARKS

What have we learned about GUT monopoles? (1) They are exceedingly interesting objects, which, if they exist, must be relics of the earliest moments of the Universe. (2) They are one of the very few predictions of GUTs that we can attempt to verify and study in our low energy environment. (3) Becuase of the glut of monopoles that should have been produced as topological defects in the very early Universe, the simplest GUTs and the standard cosmology (extrapolated back to times as early as $\simeq 10^{-34}$ s) are not compatible. This is a very important piece of information about physics at very high energies and/or the earliest moments of the universe. (4) There is no believeable prediction for the flux of relic, superheavy magnetic monopoles. (5) Based upon astrophysical considerations, we can be reasonably certain that the flux of relic monopoles is small. Since it is not obligatory that monopoles catalyze nucleon decay at a prodigious rate, a firm upper limit to the flux is provided by the Parker bound[43]; $\langle F_M \rangle \lesssim 10^{-15}$ cm^{-2} sr^{-1} s^{-1}. Note, this is not a predicted flux, it is only a firm upper bound to the flux. It is very likely that flux has to be even smaller, say $\lesssim 10^{-18}$ cm^{-2} sr^{-1} s^{-1} or even 10^{-22} cm^{-2} sr^{-1} s^{-1}. (6) There is every reason to believe that typical monopoles are moving with velocities (relative to us) of at least a few $\times 10^{-3}$ c. [Although it is possible that the largest contribution to the local monopole flux is due to a cloud of monopoles orbiting the sun with velocities $\simeq (1-2) \times 10^{-4}$ c, I think that it is very unlikely.[44-45]]

Based upon the above (not unbiased) 'list of facts', I believe that when designing a monopole detector, the monopole hunter/huntress must give highest consideration to building a detector which is sensitive to a monopole flux at least as small as 10^{-15} cm^{-2} sr^{-1} s^{-1}, for monopole velocities $\simeq 10^{-3\pm 0.5}$ c. The risks involved in monopole hunting are very great, but the potential payoff is even greater!

I apologize for any omissions I may be guilty of in this mini-review of monopoles and the standard cosmology. More complete reviews can be found in Refs. 36 and 46. This work was supported by the DoE at Fermilab and Chicago (AC0280ER10773A004), by NASA at Fermilab, and by an Alfred P. Sloan Fellowship.

REFERENCES

1. E.W. Kolb and M.S. Turner, Ann. Rev. Nucl. Part. Sci. 33, 645 (1983).
2. E.N. Parker, Astrophys. J. 160, 383 (1970).
3. G. Lazarides, Q. Shafi, and T. Walsh, Phys. Lett. 100B, 21 (1981).
4. M.S. Turner, E.N. Parker, and T.J. Bogdan, Phys. Rev. D26, 1296 (1982).
5. Y. Rephaeli and M.S. Turner, Phys. Lett. 121B, 115 (1983).
6. E.W. Kolb, S. Colgate, and J. Harvey, Phys. Rev. Lett. 49,

1373 (1982).

7. S. Dimopoulos, J. Preskill, and F. Wilczek, Phys. Lett. 119B, 320 (1982).
8. K. Freese, M.S. Turner, and D.N. Schramm, Phys. Rev. Lett. 51, 1625 (1983).
9. F. Bais, J. Ellis, D.V. Nanopoulos, and K.A. Olive, Nucl. Phys. B219, 189 (1983).
10. T. Walsh, in Proceedings of the XXIIth Intl. Conference on High Energy Physics (Paris, 1982).
11. E.W. Kolb and M.S. Turner, Astrophys. J., in press (1984).
12. For a more detailed description of the standard cosmology, see, e.g., S. Weinberg, Gravtiation and Cosmology (Wiley: NY, 1972), chapter 15.
13. J. Yang, M.S. Turner, G. Steigman, D.N. Schramm, and K.A. Olive, Astrophys. J., in press (1984).
14. P.A.M. Dirac, Proc. Roy. Soc. (London) A133, 60 (1931).
15. G. 't Hooft, Nucl. Phys. B79, 276 (1974).
16. A.M. Polyakov, JETP Lett. 20, 194 (1974).
17. T.W.B. Kibble, J. Phys. A9, 1387 (1976).
18. J. Preskill, Phys. Rev. Lett. 43, 1365 (1979).
19. Ya.B. Zel'dovich and M. Yu Khlopov, Phys. Lett. 79B, 239 (1978).
20. M.B. Einhorn, D.L. Stein, and D. Toussaint, Phys. Rev. D21, 3295 (1980).
21. A.H. Guth and E. Weinberg, Nucl. Phys. B212, 321 (1983).
22. J.A. Harvey and E.W. Kolb, Phys. Rev. D24, 2090 (1981).
23. T. Goldman, E.W. Kolb, and D. Toussaint, Phys. Rev. D23, 867 (1981).
24. J. Fry, Astrophys. J. 246, L93 (1981).
25. J. Fry and G. Fuller, Univ. of Chicago preprint (1983).
26. D.A. Dicus, D.N. Page, and V.L. Teplitz, Phys. Rev. D26, 1306 (1982).
27. F. Bais and S. Rudaz, Nucl. Phys. B170, 507 (1980).
28. A. Linde, Phys. Lett. 96B, 293 (1980).
29. G. Lazarides and Q. Shafi, Phys. Lett. 94B, 149 (1980).
30. P. Langacker and S.-Y. Pi, Phys. Rev. Lett. 45, 1 (1980).
31. E. Weinberg, Phys. Lett. 126B, 441 (1983).
32. A. Vilenkin, Tufts Univ. preprint (1983).
33. K. Lee and E. Weinberg, Columbia Univ. preprint (1984).
34. M.S. Turner, Phys. Lett. 115B, 95 (1982).
35. G. Lazarides, Q. Shafi, and W.P. Trower, Phys. Rev. Lett. 49, 1756 (1982).
36. J. Preskill, in The Very Early Universe, eds. G.W. Gibbons, S.W. Hawking, S. Siklos (Cambridge Univ. Press, Cambridge 1983).
37. W. Collins and M.S. Turner, Phys. Rev. D, in press (1984).
38. A. Goldhaber and A. Guth, in preparation (1984).
39. J.P. Vallee, Astrophys. Lett. 23, 85 (1983).
40. K. Freese, J. Frieman, and M.S. Turner, Univ. of Chicago preprint (1984).

41. M.S. Turner, Nature 302, 804 (1983).
42. M.S. Turner, Nature 306, 161 (1983).
43. The Parker bound can be evaded if monopoles play a role in the generation of the galactic magnetic field (e.g. by magnetic plasma oscillations); see Ref. 4; J. Arons and R. Blandford, Phys. Rev. Lett. 50, 544 (1983); or E. Salpeter, S. Shapiro, and I. Wasserman, Phys. Rev. 49, 1114 (1982). In order for this mechanism to work, the phase velocity of the plasma oscillations must be greater than the gravitational velocity dispersion of the monopoles $\simeq 10^{-3}$ c (otherwise Landau damping will rapidly damp the oscillations). This results in a lower bound to the monopole flux: $\langle F_M \rangle \gtrsim 1/4\, m_M v_{grav}^3 (gl)^{-2} \simeq 10^{-12} (m_M/10^{16}\text{ GeV})\text{cm}^{-2}\text{ sr}^{-1}\text{ s}^{-1}$; here $g \simeq 69e$ is the monopole's magnetic charge, $l \simeq 300$ pc is the characteristic length scale of the galactic B-field, and $v_{grav} \simeq 10^{-3}$ c. For $m_M \gtrsim 10^{18}$ GeV the monopole density required is inconsistent with the mass density density of the galaxy (see Fig. 1). I believe that this scenario is very unlikely; in addition, it is becoming observationally untenable because of the large monopole flux that it predicts.
44. S. Dimopoulos, S. Glashow, E. Purcell, and F. Wilczek, Nature 298, 824 (1982).
45. K. Freese and M.S. Turner, Phys. Lett 123B, 293 (1983).
46. P. Langacker, Phys. Rep. 72, 185 (1981).