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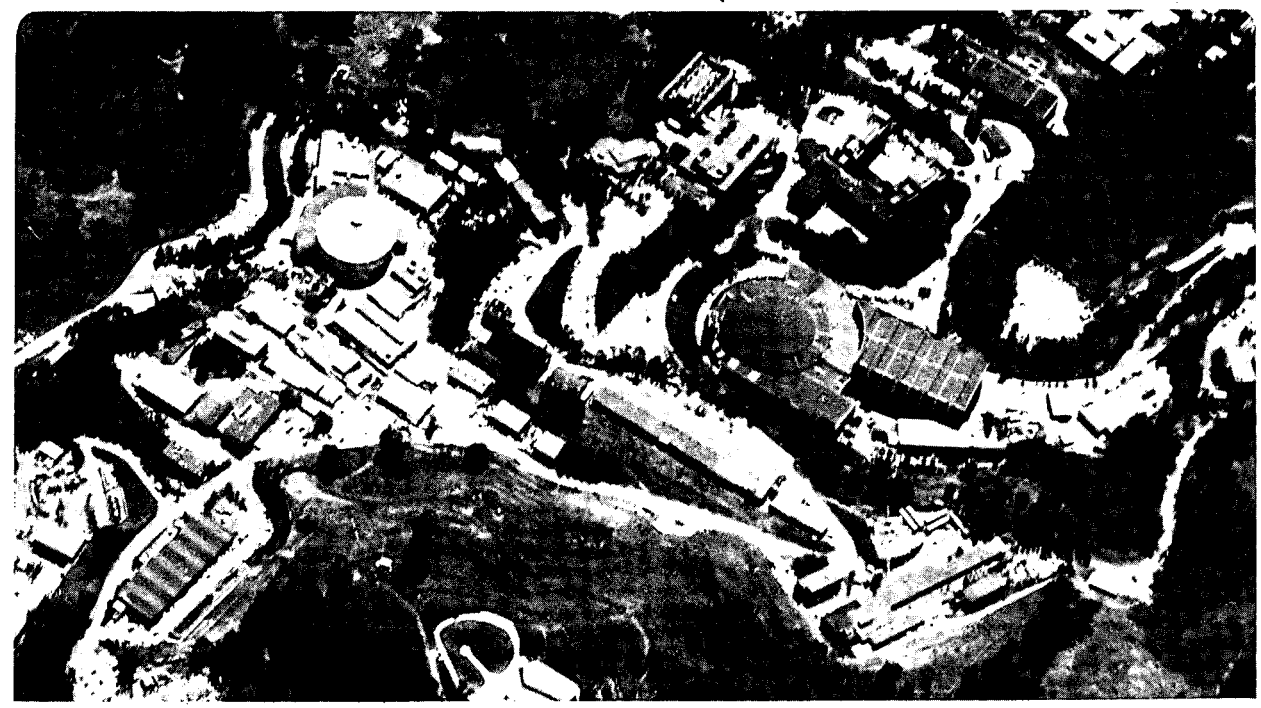
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Cosmic Relics from the Big Bang

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Cosmic Relics From The Big Bang.* †

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Abstract

A brief introduction to the big bang picture of the early universe is given. Dark matter is discussed; particularly its implications for elementary particle physics. A classification scheme for dark matter relics is given.

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(I) The Importance Of Cosmic Relics.

It is probable that the next leap forward in our understanding of particle physics will come from accelerator experiments. There are many possibilities: unexpected decay modes of K, D, B, μ, τ, Z particles and the discovery of new particles at high energies are two clear examples. Accelerator physics is crucial in unravelling the origin of electroweak symmetry breakdown. It will also shed light on flavor physics: at the very least we can learn how fermion masses are described at the TeV scale.

It is also quite possible that the next major advance in particle physics will come from astrophysics and cosmology. Most astrophysics and cosmology is done for other reasons: the questions being addressed have their own intrinsic worth. What is the nature of such objects as supernova and quasars? How did the observed mass distribution of the universe come about? What triggered clustering into galaxies, why are there so many varieties of galaxies, why do they have the size they do, and why do they themselves form clusters of galaxies? For a particle physicist perhaps the most exciting thing about astrophysics and cosmology is that there seems to be an endless succession of interesting unanswered questions. However, in these lectures I want to take a more limited viewpoint; what can we hope to learn about particle physics from astrophysics?

Infact this still leaves a wealth of possibilities open. We can certainly use our understanding of various astrophysical objects to place limits on new particles. For example, scalar particles with a mass of less than a keV and long mean free paths could be emitted from the entire volume of a star, and not just form its surface, so the existence of these particles is severely constrained. Perhaps

the best example is the supernova which went off in a nearby galaxy last year, SN1987A[1]. Although the mass limit or ν_e from this event turned out to be remarkably close to that from laboratory experiments, we did learn a great deal about particle physics from SN1987A. We learnt about other properties of ν_e: its lifetime, electric charge, magnetic moment, mixing and right handed currents. Perhaps we even learnt more about ν_r than we did about ν_e. Since we are now sure of the size of ν emission from a supernova we now know that a supernova per century in our galaxy is populating the galaxy with ν_r. If ν_r is heavy it could decay giving γ or $e^{\pm}$; in either case very stringent lifetime limits can be placed from observational limits on X and γ-ray backgrounds. We now have great confidence in these decade old limits [2].

I will reduce my scope again, and concentrate on the question of dark matter and its implications for particle physics. There is a great deal of evidence that there is much more to the universe than meets the eye. I will not discuss this evidence; there are now several books on the subject [3,4]. It is worth looking at the evidence and thinking it over for yourself. A new experimental field is opening, that of searching for dark matter [5], and we should be sure of its foundation. A typical piece of evidence has the following form: a system will be observationally analyzed to determine the mass and velocities of its constituents. One then asks whether this is a gravitationally stable system or whether more, unobserved mass is needed to stabilize it. For a great many systems; for example, stars in the local neighborhood of our galaxy, hydrogen clouds in our and other spiral galaxies, hot gas in elliptic galaxies and even for galaxies in the whole universe, it does seem that a great deal of extra mass is required. Since we have not detected this matter by means other than this gravitational dynamics, we call it dark matter.

You might guess that there is virtually no constraint on the nature of dark matter: "if we cannot see it, surely it could be anything." As with most statements containing the word "surely" it is completely false. We have three very powerful constraints which restrict the nature of dark matter:

1. We know the location of the dark matter. (Of course there could also be dark matter in locations other than those we have studied.)

2. We know that it is dark. This is especially important for the dark matter in the local neighborhood. If the dark matter is composed of particles of mass m it is raining down on us with a flux of $\simeq 10^7/(m/GeV)cm^{-2}s^{-1}$ and we just cannot see it.

3. Dark matter should result from a reasonable big bang cosmology. If you start the big bang off with a given set of particles with given interactions they typically annihilate and do not survive until today. Requiring survival with the observed abundance is a very powerful constraint on any relic object.

To implement the third constraint, it is necessary to have an understanding of the hot big bang model of the early universe. In the next section we discuss this picture, which emerges uniquely from three cornerstones: the isotropy of the 3^0K microwave background radiation, the general theory of relativity and the $SU(3) \times SU(2) \times U(1)$ gauge interactions of the elementary particles. In Section III I discuss three general points which have to do with dark matter. I give a few remarks on the experimental results, I discuss whether the dark matter could be baryonic or whether it requires an extension of particle physics to include exotic stable objects, and finally I consider inflation and its implications for dark matter. In Section IV I introduce a classification scheme for dark matter candidates and give examples and experimental signatures.

Why have I chosen to orient these lecture at the question of cosmic relics? Cosmic relics, both visible baryonic and dark non-baryonic, are the best evidence that we have for particle physics beyond the standard model.

II) A Brief Introduction To The Big Bang

The purpose of this section is to present the framework of the big bang cosmology in a simple and brief way. Many important details will be omitted and can be found together with references, elsewhere [6,7,8]. Emphasis will be on ideas rather than formalism.

A simple interpretation of the Hubble red-shift law for distant galaxy recessional velocities and of the isotropy of the 3^0K microwave radiation is that the present universe is expanding and has evolved from an early era of a hot expanding homogeneous and isotropic plasma. This system can be described at

any era by the plasma temperature $T(t)$, its pressure $p = p(\rho)$, and the chemical potentials for species i $\mu_i(T)$. These are determined from a knowledge of the constituents of the plasma and their interactions. The expansion itself is described in terms of the Robertson–Walker scale factor $R(t)$ which appears in the metric:

$$d\tau^2 = dt^2 - R(t)^2 \left(\frac{dr^2}{1 - kr^2} + d\Omega^2 \right) \tag{2.1}$$

$d\Omega$ is the usual element of solid angle, t the proper time at any location of fixed r, and r is a dimensionless coordinate. The proper distance between fluid elements at r_A and r_B at time t is given by

$$d(t) = \int_{r_A}^{r_B} \frac{R(t)dr}{\sqrt{1 - kr^2}}. \tag{2.2}$$

Hubbles law, $\dot{d} = Hd$, follows directly from this, with $H = \dot{R}/R$ being the spatially constant but time dependent Hubble parameter. I will scale the coordinate r so that the dimensionless constant k either vanishes or has unit magnitude. If $k = +1$ the universe is closed, if $k = -1$ it is open, while if $k = 0$ it is critical. Although k is crucial for the future behavior of the universe it is frequently unimportant during early times and can be set to zero.

The present value of the Hubble parameter has been measured to be

$$H_0 = 100h \ km \ s^{-1} Mpc^{-1} \tag{2.3}$$

with $1/2 \lesssim h \lesssim 1$. In many cosmologies H_0^{-1} sets the scale for the age of the universe. For example suppose the expansion of the universe is given by some power law $R = R_0(t/t_0)^n$, then $H = n/t$ and $t_0 = nH_0^{-1} \simeq nh^{-1}10^{10}$yr.

The simple picture of the expanding universe as the surface of an inflating balloon is helpful. Commoving coordinates are fixed to the surface of the balloon. A photon at time t_A with wavelength λ_A will have a stretched wavelength at some later time t_B given by $\lambda_B/\lambda_A = R(t_B)/R(t_A)$. The red-shift of a photon emitted at t_A and received at t_B is defined to be $(\lambda_B - \lambda_A)/\lambda_A$ and for $t_B >> t_A$ this just becomes $R(t_B)/R(t_A)$.

The dynamics of the expansion is given by Einsteins field equations for general relativity applied to the metric of equation (2.1)

$$\left(\frac{\dot{R}}{R} \right)^2 = \frac{8\pi G}{3} \rho - \frac{k}{R^2} \tag{2.4}$$

where G is the Newtonian gravitational constant, and ρ is the total energy density at any point in the homogeneous fluid at time t. A useful mnemonic for this equation is inspired by Newtonian ideas as depicted in Figure 1.
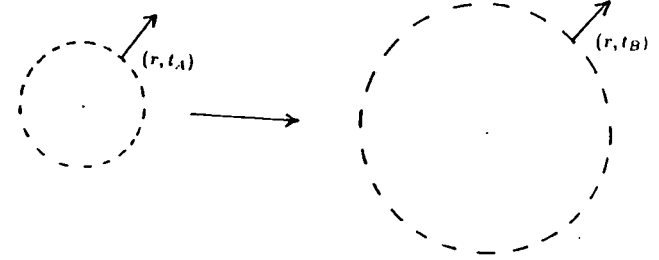
Figure 1

Newtonian mnemonic for interpretation of $\dot{R}/R$.

Imagine the expansion of a small spherical commoving region with a unit test mass at coordinate radius $r << 1$. The sum of its kinetic and potential energies $\dot{R}^2/2 - GM/R$, where M is the mass enclosed in the unit sphere, is just $-k/2$. Thus if $k = -1$ the total energy is positive and the universe is open. The Newtonian picture is not correct for the whole system. However, equations (2.4) is correct as it follows from general relativity. It is about the most important equation of big bang cosmology since it tells you how fast the universe is expanding at any time. If the universe is flat, $k = 0$, equation (2.4) can be solved to give $\rho \propto H^2$. This critical density today is

$$\rho_c = \frac{3}{8\pi G} H_0^2 \simeq 10^{-5} GeV cm^{-3} \tag{2.5}$$

I like to interpret the right hand side of equation (2.4) as the terms which drive the expansion. At early times we know that $G\rho >> R^{-2}$ so the expansion of the universe is driven by energy density rather than curvature. We do

not know what dominates the driving today; it could be the curvature term. If $G\rho_0$ dominates in driving the expansion today it is probably only through the invisible or dark components of ρ_0. The energy density today has various components: $\rho_0 = \rho_{VB} + \rho_{EM} + \rho_\nu + \rho_{DM} + \rho_V$ (VB is visible baryons, EM i.e. the electromagnetic content of the universe dominated by the 3^0K microwave radiation, ν refer to massless neutrinos, DM to dark matter and V to vacuum energy ie. a cosmological constant.) Defining Ω_i to be the ratio of density in any component i relative to the critical density

$$\Omega_i = \frac{\rho_i}{\rho_c} \qquad (2.6)$$

we know that $\Omega_{EM} \simeq 0(10^{-5}), \Omega_{VB} = 0(10^{-2})$ and for very light neutrinos $\Omega_\nu = 0(10^{-5})$. When we observe distant galaxies we see them as and where they were during a previous era: that when the detected photons left them. This leads to a violation of Hubble's law which is dependent on the deceleration parameter $q_0 = -\ddot{R}R/\dot{R}^2|_{t_0} = \Omega_0/2$. Experimental measurements of q_0 lead to $\Omega_0 \lesssim 2$. Hence we know from very direct observations that

$$10^{-2} \lesssim \Omega_0 \lesssim 2 \qquad (2.7)$$

If Ω_0 is larger than Ω_{VB} the difference is predominately due to the presence of dark matter.

At early times during the big bang the temperature was high so that particle number densities in the plasma were large enough to give reaction rates sufficient to maintain thermal equilibrium between many species of particle. For a reaction rate to be fast enough to maintain thermal equilibrium it should basically be faster than the expansion rate of the universe, that is the mean free time for interactions should be less than the age of the universe at the era under consideration. For example, for a reaction $AB \to X$ to keep particle A in thermal equilibrium at time t requires the reaction rate

$$\Gamma_{AB \to X} = n_B \langle \sigma_{AB \to X} v_{AB} \rangle > \Gamma_{exp}(t) = H(t) = \frac{\dot{R}}{R} \qquad (2.8)$$

where $R_{A,B}(t)$ are particle member densities at time t and $\langle \sigma_{AB \to X} v_{AB} \rangle$ is the thermally averaged cross-section times relative speed for this process at temperature $T(t)$. In thermal equilibrium the particle number densities are given

by

$$n_i^{(c)}(T) = \frac{g_i}{8\pi^2} \int \frac{d^3\bar{p}}{e^{\frac{E-\mu_i}{kT}} \pm 1} \qquad (2.9)$$

where g_i is the number of spin states of particle $i, \mu_i(T)$ is the chemical potential and ± 1 refers to fermions and bosons. The particle mass m_i enters via $E^2 = \bar{p}^2 + m_i^2$. Useful order of magnitude approximations are

$$n_i^{(c)}(T >> m_i) \sim T^3 \qquad (a)$$

$$n_i^{(c)}(T << m_i) \sim (m_i T)^{3/2} e^{(\mu_i - m_i)/T} \qquad (b)(2.10)$$

In the standard hot big bang model ρ is dominated by contributions from relativistic species (with $T >> m_i$, so $n_i \sim T^3$) for virtually all times for which there is a plasma which is thermally coupled; that is for all temperatures above about 1 eV. Although this is not the case for non-standard cosmologies with heavy long-lived exotic particles or with periods of inflation, the radiation dominated era is a very important one. During this era, since $n \sim T^3$ and $< E_i > \sim T$, we have $\rho \sim T^4$, so that Einstein's equation has the form

$$\dot{R}/R \sim (G\rho)^{1/2} \sim T^2/M_p \qquad (2.11)$$

When the universe expands it does work and hence cools; there must be a relationship between $T(t)$ and $R(t)$. During the radiation dominated era the relationship is that RT is constant. This is equivalent to the entropy in a comoving volume (one which grows as R^3) being constant. The gas is expanding adiabaticly under most circumstances, so that $RT =$ constant holds more generally, and is violated only when there is a mechanism which creates entropy. Using $\dot{R}/R = -\dot{T}/T$ together with (2.11) allows for a solution for $\dot{T}(t)$ during the radiation dominated era: $t \sim M_p/T^2$. The constant of proportionality does have a weak dependence on T. Near an MeV:

$$\frac{t}{sec} \simeq \left(\frac{MeV}{T}\right)^2 \qquad (2.12)$$

How does this hot plasma at $T > eV$ evolve into the universe we see today? At first sight they seem very different: apart from local hot-spots, we see a very cold universe with few particle interactions, also it is grossly inhomogeneous on

all scales up to at least 50 Mpc. However, we have come to realize that it is perfectly reasonable that the plasma of the hot big bang should evolve into a cold, non-interacting, inhomogeous universe. As the temperature of the plasma drops all the particle reaction rates $\Gamma_R(T) \sim n < \sigma v >$ fall much more rapidly than does the expansion rate $\Gamma_{exp}(T)$, so that for each reaction there is some critical temperature beneath which it is "frozen out". Furthermore, once the electromagnetic scattering processes freeze out, there is no longer any pressure to prevent mass perturbations in the plasma from undergoing gravitational growth. Although we are far from a complete picture of the resulting clustering, it can only stop once gravitationally stable systems, such as galaxies and stars, are formed.

There is a very basic question about the evolution from the hot plasma to the observed cold, inhomogeneous universe that we must address. What determines the abundance of all the stable fundamental particles in the universe today? Is it reasonable that the universe we see should contain p, e, γ in the observed ratios? If we introduce exotic stable particles into theories of particle physics can we calculate their present abundance in the universe?

The calculation of these abundances is very simple [9]. Consider a stable particle species i. As long as a process which changes the number of i particles is in thermal equilibrium, $n_i(T)$ will be given by equation (2.9). Suppose T_{if} is the freezeout temperature of the last such reaction to be in thermal equilibrium. At lower temperatures since the i particles do not decay and assuming they are not produced (for example by the decay of some other species) we have $n_i(t) \simeq (R(t)/R(t_f))^3 n_i^{(e)}(t_f)$ where t_f is the time corresponding to the temperature T_{if}. After freezeout the remaining i particles are just diluted by the volume expansion. Since $R \sim T^{-1}$ it is convenient to consider "reduced" number densities $f_i(t) = n_i/T^3$ which become independent of t after freezeout. Note that for a particle which is still relativistic at freezeout $(m_i < T_{if})$, $n_i^{(e)}(T_f)$ is given by (2.10a) so that the relic abundance $f_i \simeq 1$. However, if the particle was non-relativistic at freezeout $n_i^{(e)}(T_f)$ is given by (2.10b) so that $f_i \simeq \left(\frac{m_i}{T}\right)^{3/2} \exp\left((\mu_i - m_i)/T_{if}\right)$ which can reflect an enormous Boltzman suppression at freezeout for small μ_i and $m_i >> T_{if}$.

The above relic abundances are approximate since freezeout is treated as a

sudden process, which it is not. The physical ideas are correct however as can be seen by numerical integration of the rate equation. For example, consider a particle x for which the last x number changing process to freeze out is $x\bar{x} \to \cdots$. Ignoring a possible chemical potential the rate equation is

$$\frac{dn}{dt} = -\frac{3\dot{R}}{R}n - < \sigma_A v >_{x\bar{x}-\cdots} (n^2 - n^{(e)^2})$$ (2.13)

where $n_x = n_{\bar{x}} = n$. This has been numerically integrated and the freezeout behavior is shown in Figure 2. A

$f_x = \frac{n_x}{T^3}$

increasing

Z

T

m
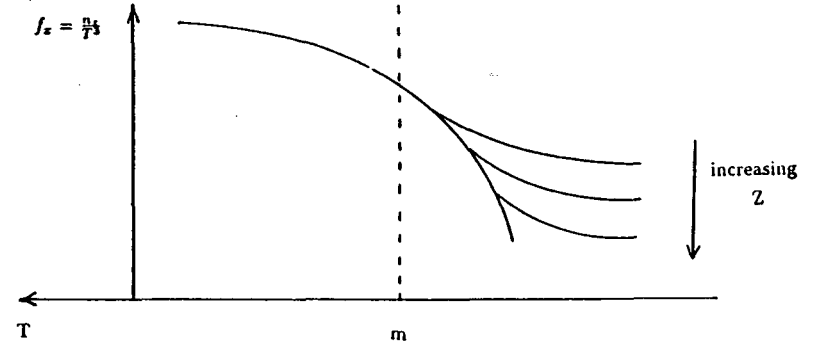
Figure 2

Schematic illustration of particle abundance freezeout behavior

useful analytic approximation for the freezeout abundance is

$$f_{xf} = \frac{n_{xf}}{T^3} \simeq \frac{\ln Z}{Z}$$ (2.14)

where $Z = m_x M_p \langle \sigma_A v \rangle$, the freezeout temperature $T_f \simeq m_x / \ln Z$ and the result is valid for large Z.

For very large Z, $f_{xf} \simeq Z^{-1}$, so that $\Omega_x \simeq (m_x f_{xf}/T_0 f_{\gamma f})\Omega_\gamma$, or

$$\Omega_x \sim \left(\frac{1}{10^3 TeV}\right)^2 \frac{1}{< \sigma_A v >}.$$ (2.15)

This result has one astounding consequence which is rarely mentioned. Even if $x\bar{x}$ annihilation proceeds via strong interactions it will be bounded by $<\sigma_A v> \lesssim \frac{1}{m_x^2}$ on dimensional grounds. Hence $\Omega_x \lesssim 2$ implies $m_x \lesssim 10^3 TeV$. It is not possible to have a stable fundamental particle with mass larger than $10^3 TeV$. This is encouraging, if the dark matter is composed of fundamental particles it is not possible to push their mass arbitrarily high, and consequently their number density and flux at the earth cannot be made arbitrarily small. This bound could only be avoided by having a phase transition give a mass to x after it has already frozen out at a low abundance.

What are the relic abundances of the four known particles (ν, γ, e, p) which we believe to be stable? The photons are massless and have relativistic freezeout $f_\gamma = 0(1)$. Photons are the only relics which were relativistic at freezeout which we have observed. They play a crucial role in determining the age of the universe. We believe the observed $3^0 K$ microwave background radiation is the photon relic of the hot big bang. These photons are no longer coupled to any plasma, but they maintain a thermal distribution with wavelength being stretched by the universal expansion: $\lambda \propto R$, so the effective temperature which describes the distribution $T \propto R^{-1}$. Knowing $T(t)$ from Einstein's equation then gives the age of the universe t_0 from the observation of $T_0 = 3^0 K$. Of the neutrinos of the standard model, ν_e and ν_μ are known to be sufficiently light that they also were relativistic at decoupling: $f_\nu \sim 1$. If they are stable and massless they contribute roughly the same to ρ_0 as do photons: $\Omega_\nu(m_\nu = 0) = 0(10^{-5})$. If their mass were $10^5 T_0 \simeq 30 eV$ they would give $\Omega_\nu(m_\nu = 30eV) = 0(1)$. This could also be true for a light tau neutrino.

If m_{ν_τ} is above 1 MeV, ν_τ would freezeout non-relativistically. For the non-relativistic freezeout of a heavy neutrino $<\sigma_A v> \sim G_F^2 m_\nu^2$ and equation (2.15) gives

$$\Omega_\nu(m_\nu) \sim \left(\frac{m_\nu}{10 GeV}\right)^2 \qquad (2.16)$$

(Since we have dropped factors of 4π equations (2.15) and (2.16) are not numerically accurate). This is excluded for a stable heavy ν_τ. If 1 MeV $< m_{\nu_\tau} < 35$ MeV, then ν_τ must be unstable. In this case the decay products typically lead to other astrophysical or cosmological problems: I expect m_{ν_τ} to be less than

about 30eV. Evidence to the contrary would have very exciting cosmological implications. Equation (2.16) applies to any neutral fermion whose dominant annihilation occurs via W and Z exchange. Such a particle with mass in the range 1-10 GeV is a good candidate for the dark matter.

At temperatures above the QCD phase transition u, d and s quarks are in thermal equilibrium with the eight gluons, and these strongly interacting particles are all relativistic. For simplicity assume that by $T = 50$ MeV the phase transition is completed and the baryon number is carried predominantly by p and n. (The picture could be much more complicated; supercooling could take place or baryon number could get trapped into quark nuggets for example.) Reactions such as $p\bar{p} \rightarrow n(\pi)$ rapidly thermalize the baryon distributions. If the chemical potential for baryon number vanishes, $\mu_B = 0$, then the freezeout abundance of p and $\bar{p}$ is given by equations (2.14), and Ω_B is given by (2.15). Since $<\sigma_A v> \sim m_p^{-2}$ this gives $\Omega_B \sim 10^{-12}$. We conclude that the standard big bang scenario requires $\mu_B \neq 0$, i.e. it requires a cosmological baryon excess to be present by $T = 50$ MeV. This conclusion is not altered by more complicated assumptions about the nature of the QCD phase transition. You might argue that it was obvious that we would need this: we do not see any evidence of primordial anti-matter anywhere in the solar system or indeed anywhere in the cluster of galaxies of which our Milky Way is a member. The only way that the entire universe could be baryon symmetric is if there were enormous domains, some baryonic and some anti-baryonic. This domain structure must have existed at $T \sim 50$ MeV to prevent over-annihilation of p with $\bar{p}$. Although an era of inflation could produce such enormous domains it has not been possible to write down a complete cosmology incorporating such a scheme.

Charge neutrality of the universe implies that the cosmic asymmetry in electrons is equal to that in protons, at least for a closed universe. This does not necessarily mean that the chemical potential for lepton number $\mu_L = \mu_B$, since there could be additional lepton asymmetries carried by neutrinos.

To obtain $\Omega_B = 0(10^{-2})$ it is necessary, just before the QCD phase transition, to have a quark asymmetry $(n_q - n_{\bar{q}})/n_q = 0(10^{-9})$. The absence of antimatter today would lead us to expect a non-zero cosmic baryon asymmetry. The importance of the standard big bang framework is that it allows as to cal-

culate how big such a asymmetry should be. The asymmetry today is enormous and obvious because essentially all anti-protons come across a proton to annihilate. However early on in the big bang it would not have been very obvious, it was a one part in a billion effect. Although small it is of crucial importance: it is non-zero and it must have come from somewhere. Assuming that it is not just a randomly adjusted initial condition to the universe, it must have been generated during an era of the universe when CP and $\dot{B}$ violating processes occurred in a non-thermal equilibrium environment. This is fascinating because it implies particle physics beyond the standard model. It also implies additional phase transitions at some early era, although the nature of the phase transition (inflation, gauge symmetry breaking ..) is pure speculation at the moment.

I will end this section with a brief summary of the main events in the hot big bang cosmology as inspired by the $SU(3) \times SU(2) \times U(1)$ model of particle gauge interactions. These events are shown in Figure 3. This model should be good up to a few hundred GeV so this is where we begin. It is likely, but not necessary, that the cosmic baryon asymmetry exists even at this high temperature. At a temperature near 250 GeV there is a phase transition at which the W, Z bosons acquire a mass as well as all the quarks and charged leptons. If the Higgs boson is light considerable supercooling (but insignificant inflation) is possible at this phase transition. I know of no feature or attribute of this phase transition which leads to observable cosmological consequences today. As the temperature drops below $M_Z, M_W, m_b, m_\tau,$ and m_c these particles are depleted by annihilation until they freezeout. Any relic abundance rapidly decays, again leaving no observable footprints. After the QCD phase transition essentially every anti-baryon annihilates with a baryon so the baryon excess now becomes an important component to the plasma rather than a miniscule asymmetry. The QCD phase transition if it is first order may lead to density inhomogeneities which could effect primordial abundance of the light elements.

At the MeV era many important events take place. This era, like the previous ones, has the expansion rate $\dot{R}/R$ driven by relativistic radiation energy density, so that the time-temperature relation is as given in equation (2.12). Since $m_n - m_p = O(MeV)$ neutron decay becomes important beneath an MeV in reducing the neutron to proton number density ratio. Since nuclear bind-

ing energies are $O(MeV)$ it also becomes possible to form nuclei which are not immediately photo-dissociated by the plasma. Most of the neutrons which escape decay are processed into ^4He nuclei, with trace quantities ending up in $^2H, ^3He,$ and ^7Li. The abundance of these nuclei predicted by nucleosynthesis in the standard hot big bang are in good agreement with primordial abundance inferred from a variety of observations. This is a very important success; there is direct experimental support for the big bang model back to times of $O(1 sec)$. The success has important consequences for deviations from the standard big bang and is especially important to the issue of whether dark matter could be baryonic, as we will discuss in the next section.
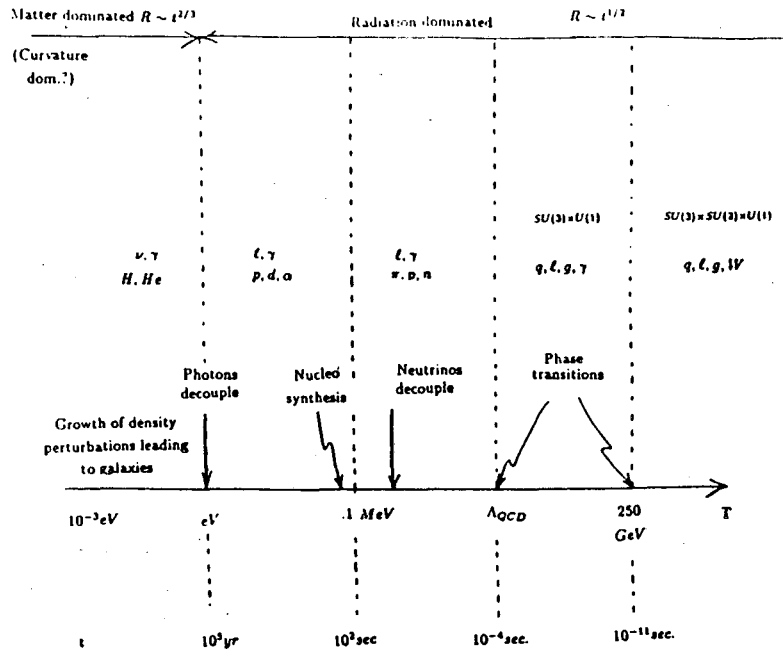
Figure 3

Important events in the hot big bang cosmology for $T < 1$ TeV.

At $O(MeV)$ the light neutrinos decouple and $e^+e^- \to \gamma\gamma$ depletes the charged leptons down to essentially just the electron excess. Hence the interacting plasma now contains electrons, protons and photons in the ration $1 : 1 : O(10^9)$, and also contains the heavy nuclei. This dilute plasma continues to cool until $T = O(eV)$. At this point $ep \to H\gamma$ takes place as the reverse photodisintegration process freezes out. Once neutral hydrogen is formed the plasma ceases to be interacting, the photons now free-stream and are red-shifted by a factor of $\simeq 3000$ until the present era.

At all temperatures smaller than $O(10MeV)$ we have $\rho_B = n_B m_B \sim 10^{-9}T^3 m_B$ and $\rho_\gamma = n_\gamma < E_\gamma > \sim T^4$. At $T = 1MeV, \rho_\gamma >> \rho_B$, however $\rho_\gamma/\rho_B \propto T$ and drops until $\rho_\gamma = \rho_B$ when $T \sim 10^{-9}m_B = 1eV$. Beneath $1eV$ the universe enters an era when $\dot{R}/R$ is driven by the hydrogen rest mass. Since $\sqrt{G\rho} \sim T^{3/2}$ and $RT =$ constant, Einsteins equation, (2.4), gives $R \sim t^{2/3}$ in this matter dominated era. If $k = 0$ this behavior continues until today. However if $\Omega < 1$ then a curvature dominated era with $R \sim t$ is reached.

The plasma at $1eV$ was homogeneous to a very high degree. We know this because we see the 3^0K microwave background to be isotropic once the peculiar motion of the earth has been accounted for. The photons of this background radiation have their last scatter at the era when $T = O(eV)$. Hence the inhomogeneities in the baryon distribution seen today, galaxy clusters, galaxies and stars must have evolved during this last factor of 3000 in redshift of the universe. The seeds for this inhomogeneity could be small density perturbations in baryons or dark matter present at $T = O(eV)$ or large density fluctuations produced by a late phase transitions at $T < 1eV$. Understanding the origin of this large scale structure which we see the universe to have is one of the most active areas of cosmological research. It is a field with many aspects: the clustering depends critically on the nature of the dominant energy density of the universe, and since this energy density is dark there are only speculations as to its nature. From the viewpoint of clustering only a few properties of the dark matter need be known, so it is sensible to group together those dark matter candidates which give essentially identical behavior. Elementary particle candidates can be divided into three such groups: hot, warm and cold. In the next section I introduce an alternative classification scheme for dark matter; one which is motivated more by the underlying particle physics.

III. Dark Matter

Any form of energy density whose existence is inferred solely from its gravitational effects is called dark matter. There is direct evidence for dark matter on all scales from the solar neighborhood to groups of galaxies, and indirect evidence that one component of dark matter may be fairly smoothly distributed over the entire universe [3,4,5]. The form of this dark matter is not known. It could be a gas of elementary particles, chunks of solid material, vacuum energy, topological defects such as monopoles, etc. and each of these classes has many particular examples.

There is evidence that dark matter is associated with galaxies of quite different types: dwarf, elliptic, and spirals, for example. I will make a few comments on the case of spiral galaxies [10], which is of particular interest as our own galaxy is spiral and from the viewpoint of direct detection of dark matter it is the local dark matter density which is of most importance. There are many different visible components to a spiral galaxy such as our own. There is the disk containing the spiral arms of stars; the sun is in such an arm about 15 kpc from the galactic center. There is also a spheroid component near the galactic center and hydrogen clouds, which extend beyond the visible limit of the spiral arms. These are sketched in Figure 4.
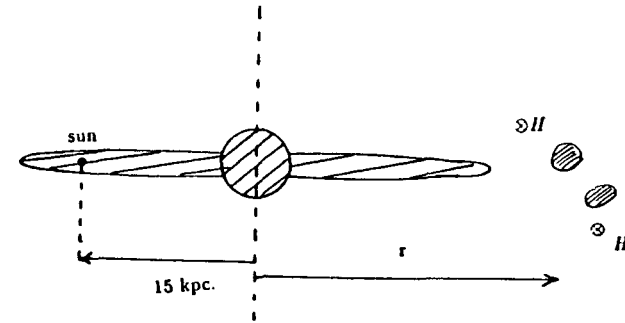

Figure 4

Schematic view of spiral galaxy.

Together they yield masses in the range of $10^{11} - 10^{12}$ times the solar mass for a typical spiral galaxy. The H clouds distant from the galactic center rotate about the symmetry axis. Since the galaxy should be in a stable gravitational configuration, and since these distant clouds provide only a small contribution to the total visible mass of the galaxy, the rotation speed should be given by $v^2/r = GM/r^2$ where M is the galactic mass. Thus one expects $v(r) \propto r^{-1/2}$. This has not been seen. Infact constant values of $v(r)$ have been observed for many spiral galaxies out to very large r (up to 100 kpc). This suggests that there is an additional component to a spiral galaxy, that of the dark halo. It should have $M(r)$ increasing as r out to very large distances, and it is frequently estimated that the halo mass is an order of magnitude larger than the visible mass.

From the viewpoint of particle physics the first important question to ask about dark matter is: is it baryonic? Examples of baryonic dark matter would be cool stars with low luminosities or perhaps planetary sided lumps of cold solid material. If this were the dark matter it would be interesting for particle physics, but certainly not revolutionary. It would simply mean that the baryonic

freezeout abundance,which is related to the cosmic baryon asymmetry, is larger than previously thought. To ignore the baryonic option simply because we do not have an understanding of how such objects are formed is a mistake: formation of all types of stars and galaxies are, to varying degrees, not understood.

One constraint on Ω_B comes from big bang nucleosynthesis [11]. As the baryonic density at the MeV era is increased so the nucleosynthesis reaction rates are increased. This depletes the low mass nuclei ($^2H,^3He$) since they are more fully burnt to the higher mass nuclei ($^4He,^7Li$) whose abundances consequently increase. The preferred range of Ω_B from a comparison of big bang nucleosynthesis with the observationally inferred abundances is a few percent. However, there are many uncertainties to do with the interpretation of the various observations, and consequently one cannot use this alone to rule out the possibility that all the known dark matter ($\Omega_{Dark} \simeq .1$) is infact baryonic.

It is widely believed that at some temperature, presumably larger than the weak scale, the universe began an era when the Robertson-Walker scale factor $R(t)$ underwent a very rapid increase [12, 13]. This inflationary era would also produce very rapid cooling ($RT = const.$). This era is ended by some very non–adiabatic process which releases a stupendous amount of entropy and reheats the universe. There are several theoretical reasons as to why this quite bizarre history is attractive to many cosmologists. I will describe the way in which inflation solves the "flatness" problem.

Recall Einsteins equation for the rate of expansion of the universe; it has the form $(\dot{R}/R)^2 \sim G\rho - 1/R^2$. Direct observations today tell us that the curvature term cannot be larger than about ten times the energy density term: $1/R^2 \lesssim 10G\rho$. As we go to earlier times ρ increases first as T^3 (matter dominated era) then as T^4 (radiation dominated era) while R^{-2} increases only as T^2. Thus at very early times $G\rho$ dominates by an enormous amount. Infact at the Planck scale the initial condition required for evolution to reach the present universe is

$$\frac{1/R^2}{G\rho}\bigg|_{T=M_p} \lesssim 10^{-60}. \tag{3.1}$$

A much more natural initial condition would be for this ratio to be unity. In that case the universe will rapidly become curvature dominated: $\dot{R}/R =$

1/R or $R = t$. Furthermore with $RT = $ constant, the initial condition implies $T \simeq t^{-1}$. Since $\rho \sim T^4$ and $\rho_c \sim H^2/G \sim M_p^2 T^2$ we find that in this universe $\Omega(T) \sim T^2/M_p^2$. Thus when $T = T_0 = 3^0K$, $\Omega_0 \sim 10^{-60}$. This temperature of the universe is reached at a time $t_0 \simeq T_0^{-1} \sim 10^{-12}$ sec. Naturalness implies that open universes become cold very quickly, certainly this is not our universe. You might argue that to avoid this problem just set $k = 0$ (I took $k = -1$ in the above). This is avoiding the issue: putting $k = 0$ is the same fine tune as making the $1/R^2$ term negligible compared with $G\rho$.

The inflationary solution to this problem is sketched in Figure 5. There is a natural initial condition $G\rho \sim 1/R^2$
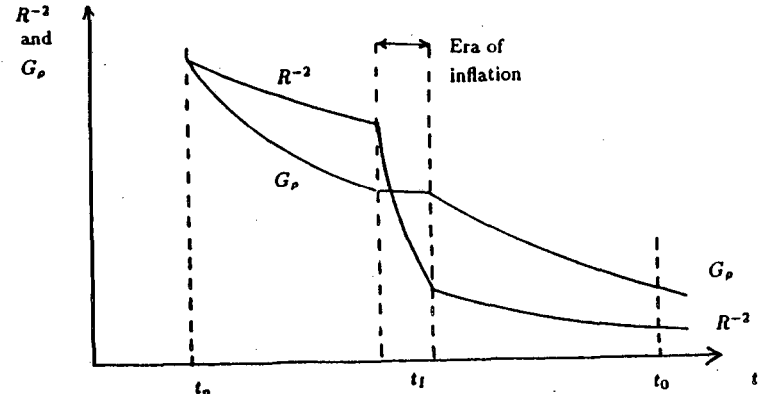

Figure 5

Inflation solves the flatness problem.

at $T \sim M_p$. A curvature dominated era then takes place, but long before the present era, the $1/R^2$ curvature term is made very small by inflation: recall that inflation rapidly increases R. You might think that since T drops rapidly ρ would also decrease catastrophically. That is not correct; during inflation ρ is dominated by vacuum energy density ρ_V which is constant. It is true that the radiation energy density $\rho_r \sim T^4$ does drop, however at the end of inflation it is replenished because ρ_V is converted to radiation energy density. After inflation $G\rho$ is much larger than R^{-2}. While $G\rho$ subsequently drops faster than R^{-2} there

is no reason to expect that we are now in the era when they are comparable, indeed this would itself constitute a fine tune. Hence inflation gives $G\rho \gg R^{-2}$ today, ie it gives $\Omega = 1$.

The visible contributions to Ω are a few per cent. Known dark matter contributions give an Ω of ten to twenty percent. In view of the theoretical motivation from inflation, it seems to be a small step to assume that further dark matter is out there and that Ω is really very close to unity. However, from the viewpoint of particle physics this is the crucial step: standard nucleosynthesis allows $\Omega_B = .1$, but $\Omega_B = 1$ really is excluded since the deuterium abundance is then three orders of magnitude too small. Inflation is tremendously exciting: it not only solves cosmological problems such as the flatness problem but it dictates considerable new particle physics beyond the standard model. There should be new exotic stable objects which contribute perhaps 90% to Ω and are not clustered on galactic scales (I'll call this the diffuse dark matter component). There must also be the particle physics which is responsible for the vacuum energy which drives the inflation.

It is also clear that we should tread very carefully: suppose somebody invents a new idea which plays the role of inflation but which does not require $\Omega = 1$. In this case there is no need for the diffuse component of dark matter. In this case there is still the issue of the clustered dark matter. It could be baryonic or exotic, and both forms are worth searching for. It is also worth stressing that the case for exotic dark matter also relies on our understanding of nucleosynthesis. Both theory and interpretation of observation have many interesting unresolved issues. An important theoretical question currently under study is whether inhomogeneities created at the QCD phase transition could have persisted to the nucleosynthesis era with sufficient size to radically alter the conventional abundances which are predicted assuming homogeneity [14]. This is not yet resolved; it seems that in a small region of parameter space $\Omega_B = 1$ might yet be consistent with the standard model.

A more radical question is whether we have identified the main era of light element abundance in the big bang. It is possible to reproduce acceptable abundances for $^2H,^3He,^4He$ and ^7Li with $\Omega_B = 1$ in schemes which have an exotic particle decaying during the keV era to produce showers in which nucleosyn-

thesis is rekindled [15]. From the viewpoint of particle physics, it is not clear that a stable exotic particle should be preferred over an unstable one. Form the viewpoint of observation it is important to know whether or not the dark matter is baryonic or exotic. Such late decaying schemes predict a higher primordial abundance of ^6Li than the standard model, and this can be searched for [16].

It is possible that the dark matter has several components. This certainly complicates issues such as galaxy formation, but it is not unreasonable: we know of many objects which arise quite naturally in gauge theories which could survive the big bang and would be dark today. I will divide these components into two classes, those which clump and those which do not.

The clumped components contribute about .1 to Ω. They are clumped on galactic scales and hence are non-relativistic. If these components are elementary particles their typical speeds will be $O(10^{-3})c$, this is the typical infall velocity into a galaxy. Faster speeds would lead to the particles escaping from the galaxy. The clumped components could be baryonic or exotic. The components which are not clumped on galactic scales could contribute up to $\Omega = 1$. Such a large, smooth contribution to Ω is certainly not baryonic. To avoid clumping it should be something like relativistic particles or vacuum energy.

If there is a dark matter halo in our galaxy composed of exotic particles, why do we not see these particles in the earth? To understand this it is useful to remember how stars form in a galaxy. A condensation of a pre-stellar gas cloud occurs only because the components (H_2 molecules for example) can sink in the gravitational potential of the cloud by dissipating their energy. For example molecular collisions excite rotational modes which give photons on de-excitation: kinetic energy is radiated away from the cloud. This dissipation is crucial for any gravitational collapse to form a stable dense object. Presumably the reason why the earth and sun do not contain enormous quantities of dark matter is that the halo dark matter particles are sufficiently weakly interacting that they are unable to dissipate this kinetic energy, they have no option but to move on bound orbits around the galaxy. This is quite reasonable. Since they are dark, these particles cannot have strong or electromagnetic interactions. Presumably their interaction strength is that of the weak interactions; they are often called WIMPS: weakly interacting massive particles.

It is a mistake to think that the scattering cross-sections for these particles is really $O(G_F^2)$. A sufficient requirement for a halo stable against dissipation is simply that the mean free time for collisions of a dark matter particle in the halo, Γ^{-1}, should be longer than the age of the galaxy, τ_G. This allows a cross-section $< \sigma\beta >$ for dark matter particles scattering from baryons as large as a millibarn, hardly a weak cross-section.

I will finish this section with a few simple estimates of event rates which could be expected if the dark matter of our halo is composed of particles of mass m_D. From rotation curves of our galaxy and other observations, the best estimate for the local density of halo dark matter is .3 GeV cm^{-3}. Thus with speeds of $10^{-3}c$ the flux expected at the earth is $\sim 10^7 cm^{-2}s^{-1}(GeV/m_D)$. Suppose that we build a detector to try and observe this enormous flux. If the dark matter particles scatter from the nuclei of the detector, of mass m_N, with some cross-section $\sigma_{S,N}$, then the event rate per kilogram of detector is

$$\left(\frac{Events}{kg\ Day}\right) \simeq \left(\frac{\sigma_{S,N}}{2.10^{-39}cm^2}\right)\left(\frac{GeV}{m_D}\right)\left(\frac{GeV}{m_N}\right). \qquad (3.2)$$

For large enough cross-sections the rates could be enormous. However, detecting the events is a considerable experimental challenge. Since the dark matter particles are non-relativistic with $\beta = 10^{-3}$, their kinetic energies are

$$T_D \simeq \left(\frac{m_D}{GeV}\right) keV \qquad (3.3)$$

and only a fraction of this will appear as nuclear recoil. The challenge is to build detectors which can measure energy depositions of $O(keV)$. Searches have already been performed with low background Ge detectors which have masses of $O(1Kg)$ and which were originally designed to search for double beta decay. The approximate excluded region in the $m_D, \sigma_{S,Ge}$ plane is shown in Figure 6.
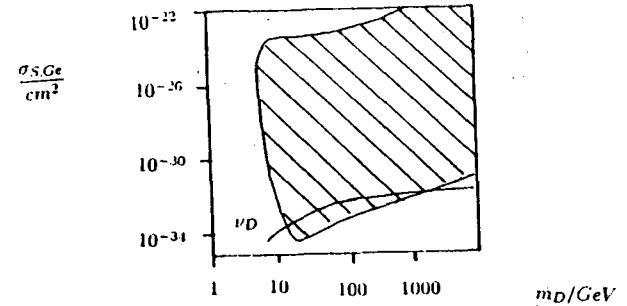
Most interesting is the

Figure 6

Excluded region in the $m_D/\sigma_{S,Ge}$ plane assuming the particle is responsible for the local dark mass. The upper edge of the shaded region corresponds to the cross-section for scattering from the rock overburden. The plot is taken from reference x and results from use of a double beta decay Ge spectrometer. The curve labelled ν_D is for a Dirac neutrino with conventional weak interactions.

region $1 < m_D < 10$ GeV, which has not yet been excluded. The reason that this region is so interesting is straightforward. The necessary annihilation cross-section σ_A for a cosmic relic to survive the big bang and contribute $\Omega \simeq .1$ today can be read from equation (2.15). For several interesting candidates, such as a Dirac neutrino, σ_A is a known function of m_D and hence one can predict m_D and it turns out to lie in this region of a few GeV. Furthermore, if the annihilation process is to ordinary matter, eg $D\bar{D} \rightarrow q\bar{q}$, then by a crossing relation it is possible to calculate the scattering cross-section from nucleus N : $\sigma_{S,N}$. Putting these values of m_D and $\sigma_{S,N}$ into (3.2) one finds that the event rates are quite large. The problem with seeing these events can be seen from equation (3.3). The lower value of m_D decreases the kinetic energy of the dark matter particle.

Future detectors will be sensitive to such low energy depositions and will explore this crucial region.

Our failure to directly detect dark matter on the earth can have a variety of explanations. Typically it is either that $\sigma_{S,N}$ is too small, m_D is so large that the event rate is too low, or m_D is so small that the signal is too feeble.

IV a Classification Scheme For Cosmic Relics

Cosmic relics can be classified according to the way in which they are produced in the big bang and the way in which they survive until today. All relics of which I am aware fall into one of six classes. For three of these classes the relics were once in thermal equilibrium and they underwent a freezeout process. Relics of the remaining three classes were produced in catastrophic events such as phase transitions and were never in thermal equilibrium. We discuss each class in turn and the classification scheme is summarized in the Table.

IV.1 Plasma Relics

Plasma relics are elementary particles which are relativistic when they decouple. That is at decoupling their number density is $n_D \sim T_D^3$, and today $n_0 \sim T_0^3$. If their mass is less than T_0 then $\rho_0 \sim T_0^4$ so $\Omega_0 \sim 10^{-5}$. The three degree microwave photon background is the best illustration of a plasma relic which is still relativistic today. If the masses for ν_i are less than T_0 they are also plasma relics which are relativistic today. Such plasma relics are not important for dark matter, unless there were $O(10^4)$ such species. This bizarre possibility is excluded since they would greatly increase ρ and therefore $\dot{R}/R$ at the time of nucleosynthesis, thus destroying the successful predictions of primordial nucleosynthesis. Hence although plasma relics which are still relativistic today will not be clustered, they cannot contribute much to Ω.

Plasma relics which have masses larger than T_0 would be non-relativistic today and would clump. Any of the three neutrinos could have such masses and could therefore make an important contributions to Ω.

IV.2 Freezeout Relic

This important case was discussed in the previous sections. Freezeout relics are particles which were once in thermal equilibrium and are non-relativistic when the reactions which change their comoving number densities freeze out. Their contribution to Ω is given by equation (2.15) for any such species x. Although the mass does not appear explicitly in Ω_x, their is almost always implicit dependence via the annihilation cross-section. Since $< \sigma_A v > \lesssim 1/m_x^2$ even for a strongly interacting particle, a freezeout relic is expected to be lighter than 10^3 TeV. For a particle with a weak annihilation cross section $< \sigma_A v > \simeq G_F^2 m_x^2$

so that $\Omega_x \simeq .1$ for halo dark matter results with $m_x \simeq O(GeV)$. Although these candidates, such as Dirac neutrinos or supersymmetric photinos are quite plausible, they are not the only possible freezeout relic.

Consider a new version of QCD, shadow QCD, which has an asymptotically free gauge coupling which gets strong at Λ' causing confinement of shadow quarks q' into shadow baryons B' which acquire mass $O(\Lambda')$. These baryons can annihilate into π' : $B'\overline{B}' \rightarrow \pi' \cdots$ with $< \sigma_A v > \sim 1/\Lambda'^2$. If $\Lambda' \simeq 300$ TeV these shadow baryons would give $\Omega \simeq .1$ even if there were no cosmic B' asymmetry.

Similarly one could imagine a world with a new unbroken $U(1)$ gauge group: shadow QED. If m'_e is the mass of the shadow electron, the lightest particle carrying shadow charge, then $e'\overline{e}' \rightarrow \gamma'\gamma'$ has $< \sigma_A v > \simeq \alpha'^2/m_{e'}^2$ so that $\Omega_{e'} \simeq .1$ arises with $m_{e'} \simeq \alpha'300$ TeV, again taking zero shadow lepton asymmetry.

IV.3 Asymmetric Relic.

Asymmetric relics are freezeout relics whose survival abundance has been greatly enhanced because of a cosmic particle anti-particle asymmetry. Protons and electrons are the best examples. It is very plausible that the dark matter is an asymmetric relic: we know native produced a cosmic asymmetry in baryon number B, so it is reasonable that it has done the same for some other quantum number. Dark matter would have survived until today for precisely the same reason that the visible matter did.

It is fun to redo the Lee–Weinberg [9] freezeout calculation including a chemical potential. In particular, while the chemical potential directly determines the abundance of the surviving major component, the minor component is annihilated way below what would have survived with zero chemical potential. None of the antiprotons seen in cosmic rays survived directly from the big bang, they were made recently in high energy collisions.

A very intriguing possibility arises if the halo dark mater is an asymmetric relic. The sun can gravitationally bind dark matter particles by scattering them into a bound orbit as they pass through the sun. Over the age of the sun significant concentrations of dark matter particles could have built up in the sun, providing they are asymmetric relics so that $x\overline{x} \rightarrow \cdots$ does not deplete the concentration. It has been found that for certain masses such bound relics could

contribute to the thermal opacity of the sun, decreasing the central temperature of the sun and decreasing the reaction rate which produces the high energy solar neutrinos thus solving the solar neutrino problem [17]. While a freezeout relic (no asymmetry) can also be trapped by the sun, $x\overline{x} \rightarrow \cdots$ will prevent buildup of a sufficient concentration to be important in changing the solar opacity. However, $x\overline{x} \rightarrow \cdots$ may itself result in high energy neutrinos (now much higher in energy than usual solar neutrinos) which would be an interesting signature for some freezeout relic candidates.

IV.4 Oscillaton Relic

Consider a scalar field $\phi(\overline{x}, t)$. If the zero temperature potential for this field is as sketched in Fig. 7.
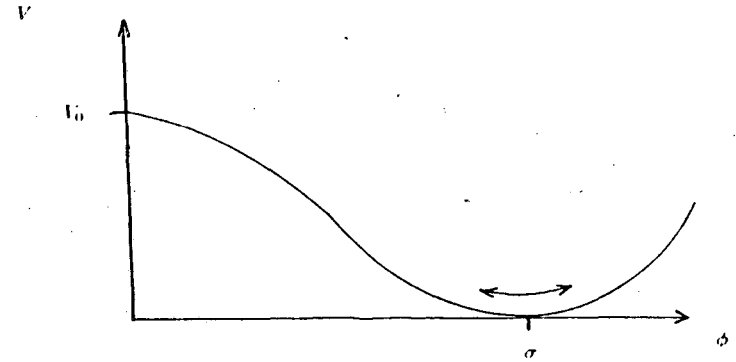

Figure 7

Typical potential for an oscillaton relic

then at some critical temperature T_0 a phase transition will occur: initially $\phi = 0$ everywhere, while at the phase transition it rolls to the minimum of the potential $\phi = \sigma$ everywhere. If ϕ has strong couplings it can radiate the energy density V_0 easily, so the equation governing the evolution of ϕ has a solution which is strongly damped. On the other hand, if ϕ has only very weak interactions the equation of motion will have small damping and the solution will be oscillatory.

In the limit that the damping can be neglected

$$\phi(\bar{x},t) = \sigma + \phi_0 e^{imt} \qquad (4.1)$$

where m is the mass of the quanta which the field ϕ creates. These oscillatons which are initiated by the phase transition would survive until today. What does this oscillaton represent physically? Because the oscillatons are the same at all spatial locations it is a mode of the field which carries no momentum. The energy density in the oscillaton represents a uniform distribution of ϕ quanta at rest; $\rho_0 = nm$. Of course, equation (4.1) is not quite correct: as the universe expands the number density of these particles gets diluted $n \sim 1/R^3$, so ϕ_0 also has a time dependence due to the expansion. However, even if these non-relativistic particles had only a small contribution to ρ at the phase transition $\rho_0(T_c) << \rho(T_c)$ they could easily dominate by today if T_c falls in the radiation dominated era. This is illustrated in Figure 8.
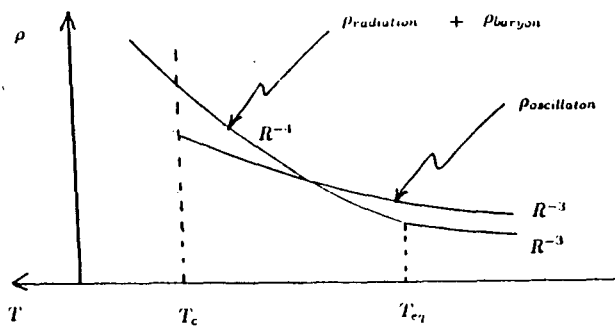
Figure 8

Temperature evolution of ρ in oscillatons and ρ in radiation and baryons. The oscillatons are produced at T_c. At T_{eq}, $\rho_{radiation}$ and ρ_{baryon} are equal.

The most well know example of an oscillaton is the axion. In the absence of QCD the axion is infact a massless Goldstone boson which results from the breaking of a global $U(1)$ symmetry, the Peccei-Quinn symmetry, at some scale f. The QCD interactions produce the potential of Fig. 7, $V_0 \sim \Lambda^4$ where Λ is the QCD parameter, and the order parameter for the symmetry breaking, σ, is f. However, oscillatons occur very frequently; they result from symmetry breaking with a weakly coupled scalar, and are much more general than the axion.

IV.5 Secondary Relics

So far we have assumed that relics are stable, or at least that their lifetimes are longer than the age of the universe. However, any of the relics considered so far could have lifetimes less than the age the universe. Their cosmologically stable decay products I will call secondary relics. They are of particular interest for obtaining $\Omega = 1$ in a smooth distribution without galactic size clumping. None of the first few classes of relic lead to dark matter today which is both relativistic, unclumped and gives $\Omega = 1$. However, a secondary relic could be relativistic today even if it came from a non-relativistic primary.

There are many examples of secondaries. The inflaton is an oscillaton. However equation (4.1) is insufficient to describe its oscillation because it is unstable so the oscillaton gets suppressed by $\exp(-\Gamma t)$. In this sense, everything is a secondary relic, we owe our existence to the decays of inflatons which reheated the universe after inflation.

As another example, consider supersymmetric theories where the lightest superpartner, which we take to be the photino, although long lived does decay via small R parity violating interactions. One example, $\tilde{\gamma} \rightarrow e^+ e^- \nu$, is illustrated in Fig. 9. This gives $e^\pm, \nu$ as secondary relics.
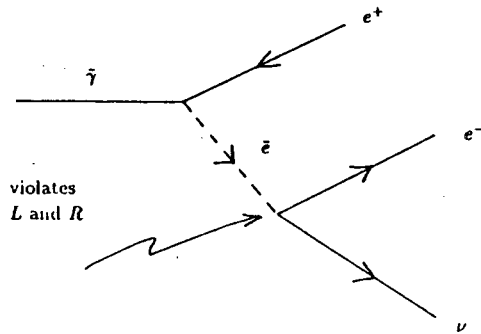

Figure 9

Photino decay via lepton and R parity violation.

In this case we can follow the evolution of the relativistic $e^\pm$ and demonstrate that they cannot contribute significantly to Ω. The $e^\pm$ lose energy rapidly by inverse Compton scattering from the background plasma photons $e\gamma \rightarrow e\gamma$. Once the $e^\pm$ are non-relativistic they clump. However, the majority of the photino rest mass has ended up in electromagnetic radiation which today would be X and γ rays. We know from background X and γ rays observations that Ω in these components are very small ($< 10^{-8}$). As usual, one finds that the big bang is a tightly constrained framework. The majority of new particle physics ideas for creating cosmic relics simply do not work. They lead to universes quite unlike our own.

IV.6 Soliton Relics

By "soliton" I mean energy which is spatially localized and which is produced at a phase transition. This is not the same as other uses of the term. Some solition relics are topologically stable defects: domain walls, strings and monopoles [18, 19]. For example consider a theory which contains a real scalar field $\phi(\vec{x}, t)$ which has a potential $V(\phi) = \lambda(\phi^2 - \sigma^2)$. This potential has two discrete degenerate minima. As shown in Fig. 10(a), if a phase transition takes place such that the field ϕ takes on a vacuum value $+\sigma$ near region A
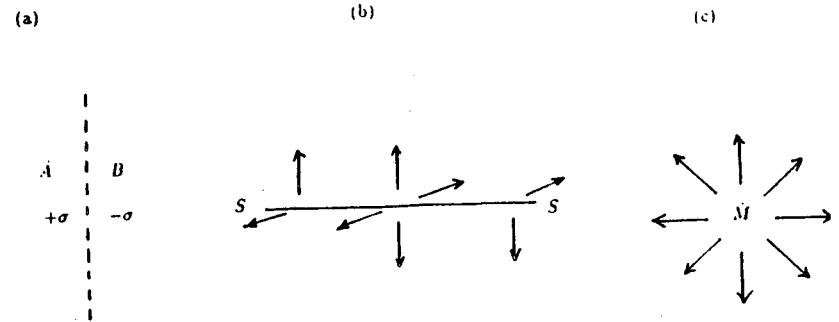

Figure 10

Topologically stable defects (a) a 2-dimensional domain wall (b)

a 1-dimensional string S (c) a point monopole M. In cases (b)and (c) the arrows represent the direction of the scalar field in internal space.

and $-\sigma$ near region B, then these regions will be separated by a domain wall at which ϕ is not at either of the minima. This domain wall contains localized field energy. If the ϕ is now made complex and the theory processes a $U(1)$ phase invariance then $V(\phi) = \lambda(\phi^*\phi - \sigma^2)^2$. In this case a pattern of vacuum field configurations result in a topological string as shown in Fig. 8(b). Infact for the energy of this field configuration to be localized on the line defect the $U(1)$ should be gauged. A monopole arises in gauge theories when a non-Abelian internal symmetry group is broken, in the monopole case the defect occurs at a point as shown in Fig. 8(c). A simple example is a theory of three real scalar fields $(\phi_1\varphi_2\varphi_3)$ which has a potential which has an $SO(3)$ invariance: $V(\phi) = \lambda(\phi_i^2 - \sigma^2)$.

The calculation of production rates for vacuum defects is not straightforward in particle collisions or in the big bang. Simple estimates for a phase transition in the big bang can be made. Suppose that at $T > T_c, \phi = 0$ in each of the above examples, while at T_c ϕ makes a transition to $|\phi| = \sigma$ everywhere except near the defects where it vanishes. The direction of ϕ is random on scales of the correlation length ζ of the phase transition. Hence, as an order of magnitude estimate the defect number density is $n_D(T_c) \simeq \zeta^{-3}$. Since ζ is certainly less than the horizon there will typically be many defects per horizon volume at T_c.

The subsequent evolution of the collection of defects can be very intricate. However, under certain simplified conditions $\rho_{DW} \sim R^{-1}$, $\rho_s \sim R^{-2}$ and $\rho_M \sim R^{-3}$ for domain walls, strings and monopoles. I know of no complete cosmologies where ρ today is dominated by domain walls or cosmic strings. Domain walls rapidly overwhelm other contributions to ρ and lead to universes quite unlike our own, unless they can be made to disappear. Cosmic strings which self-intersect and produce loops which can then disappear by gravitational radiation may not be problematic. At any era ρ_s then scales the same way with R as ρ_{TOTAL}. Today are finds $\Omega_s \lesssim \sigma/M_p$. For σ of 10^{16} GeV it is possible that the

string network may provide the inhomogenities about which galaxy clustering first occurs.

If monopoles are made at a very early phase transition their number density must be depleted by a subsequent era of inflation otherwise the universe will not evolve to the one we see. It is not possible that a monopole with magnetic charge and mass near the grand unification scale is depleted just enough to be the dark matter today. This would produce a flux of monopoles at the earth which has been experimentally excluded. It is possible that the dark matter could be a monopole which carries some other charge.

There is a second class of soliton relics. There are regions of false vacuum which have become stabilized by the presence of matter or charge. These non topological solitions I will call false vacuum nuggets. The most well known example is that of quark nuggets which could arise during the QCD phase transitions and which I describe below [20].

In Fig. 11 I sketch how the QCD phase transition would proceed cosmologically assuming that it is first order, as indicate by lattice calculations. At first small
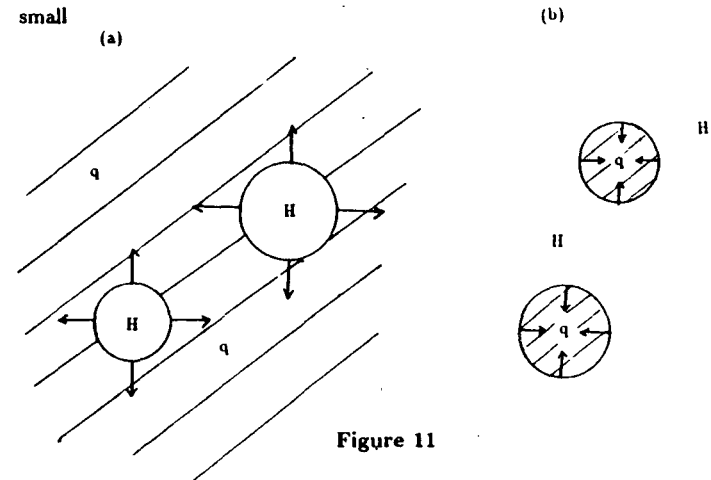

Figure 11

The cosmological QCD phase transition (a) Nucleated bubbles of hadron phase expand into the quark plasma (b) Shrinking bubbles

of quark phase.

bubbles of hadronic phase are nucleated in the previously homogenous quark gluon plasma. Since the vacuum energy of these hadronic bubbles is lower than in the quark phase the bubbles expand as shown in Fig. 11(a). The bubbles will collide and coalesce until half of space is filled by the hadronic phase. Subsequently, the picture is that regions of quark phase which are collapsing within the hadronic medium as shown in Fig. 11(b). The original quark plasma had nearly equal numbers of quarks and antiquarks, although there was a quark excess of one part per billion. However, if the critical temperature is say 100 MeV the many quarks and antiquarks must annihilate because in the hadron phase at $T_c, n_B \simeq n_{\bar{B}} \simeq (T_c, m_B)^{3/2} e^{-m_B/T_c} << T_c^3$ (see 2.10). $q\bar{q}$ annihilation can occur to $\nu\bar{\nu}$ via the Z, and since neutrinos transport energy over large distances thermal equilibrium at T_c is maintained. Now imagine following the quark excess. When the phase boundary moves into the quark fluid the quark excess tends to be swept along with the boundary. This is because for baryon number to go across the boundary energy must be found to create the baryon mass. To some degree it is therefore energetically favorable for the quark excess to remain in the quark plasma. If this effect is quite powerful then the collapse of a quark bubble shown in Fig. 11(b) will eventually be prevented by the stabilizing effect of the pressure of the quark excess inside the bubble. A quark nugget has been formed at temperature T_c.

We do not know whether such nuggets would be stable at zero temperature, they might just decay to ordinary baryons. However, it has been argued that even if quark nuggets were formed at T_c and even if they were stable at $T = 0$, they do not survive the cosmological evolution from T_c to $T = 0$, rather they evaporate [21]. While it seems that quark nuggets are not the dark matter, it is possible that some false vacuum nugget of another phase transition contributes to the dark matter.

Summary

The hot big bang model of the early universe provides a simple and elegant framework in which to study the effects of various gauge models of particle physics on cosmological issues. The standard gauge theory apparently does not lead to the universe which we observe. There is the need for baryogenesis and there is the need for a cosmic relic to be the dark matter. The cosmological description would also be more acceptable if particle physics gave rise to an era of inflation, when $R(t)$ grew very rapidly.

Most additions to the standard model which give cosmic relics do not give them with the correct abundance. The requirement of $0.1 \simeq \Omega \simeq 2$ places considerable restrictions on the interactions which generate the relics. Nevertheless an enormous number of candidate relics have been proposed. I have introduced a classification scheme to describe these cosmic relics, and it is summarized in the Table. Plasma and freezeout (including asymmetric) relics are particles which were once in thermal equilibrium with the plasma of the hot big bang. Secondary relics are particles which arise from the decay of any of these three types of primary particle relics. Oscillatons and solitons are directly associated with phase transitions in the quantum field theory, and occur in a surprising variety of forms.

As indicated in the Table, there are many ways to search for the various relics. Several searches have been done for many years and new ones with novel techniques are planned for the future. The discovery of any of these relics would be a major turning point in cosmology. It would also give us solid guidance in understanding particle physics beyond the standard model.

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Category	Charachteristics	Examples	Possible Implications	Detection
Plasma	$m < T_0$ Decoupled	γ ν_e, ν_μ?	New light Particles eg Goldstone Bosons	Difficult as low energy
Freezeout	$m > T_0$ Froze out while non-relativistic	$\nu_{M,D}$ $\tilde{\gamma}$	New physics at weak scale	Direct detection Detection of products from annihilation in sun or halo.
Asymmetric	$m > T_0$ Abundance determined by cosmological asymmetry	p, e ν_D	CP violation New quantum number	Direct detection
Oscillaton	Very weakly coupled scalar	axion	New phase transition.	Difficult as weakly coupled. Axion to γ conversion
Secondary	$X_1 \rightarrow X_2 \ldots$ $\tau(X_1) \lesssim t_0$	$\nu_H \rightarrow \nu_L \gamma$	An approximate symmetry	Other decay products of X_1 may give γ, X ray signals.
Soliton	Defects from phase transitions with localized energy.	Domain walls Strings Monopoles Nuggets	New phase transiton	Direct searches Gravitational lensing Gravitational radiation

TABLE

A classification scheme for cosmological relics.