



# Fermi National Accelerator Laboratory

FERMILAB-Conf-91/179-A  
June 1991

## TESTING THE BIG BANG: LIGHT ELEMENTS, NEUTRINOS, DARK MATTER AND LARGE-SCALE STRUCTURE \*

David N. Schramm  
The University of Chicago,  
5640 S. Ellis Avenue, Chicago, IL 60637  
and  
NASA/Fermilab Astrophysics Center,  
Fermi National Accelerator Laboratory  
Box 500, Batavia, IL 60510-0500

11-22-1991  
N-90-11-  
20328  
p40

### ABSTRACT

In this series of lectures, several experimental and observational tests of the standard cosmological model are examined. In particular, detailed discussion is presented regarding (1) nucleosynthesis, the light element abundances and neutrino counting; (2) the dark matter problems; and (3) the formation of galaxies and large-scale structure. Comments will also be made on the possible implications of the recent solar neutrino experimental results for cosmology. An appendix briefly discusses the 17 keV "thing" and the cosmological and astrophysical constraints on it.

(NASA-CP-1991-1) TESTING THE BIG BANG:  
LIGHT ELEMENTS, NEUTRINOS, DARK MATTER AND  
LARGE-SCALE STRUCTURE (Fermi National  
Accelerator Lab.) 40 p

N91-29111

CSCL 036

Unclass

03/90 0032328

\* Prepared for *Proceedings of University of Puerto Rico Second Winter School of Physics, 1991*



## Introduction

These lectures will examine several topics where the modern cosmological model is being experimentally and/or observationally examined. Some specific areas are: (1) nucleosynthesis and neutrino counting; (2) the dark matter problems; and (3) the formation of galaxies and large-scale structure. Comments will also be made on the possible implications of the recent solar neutrino experimental results for cosmology. An appendix will discuss the  $17\text{keV}$  “thing” and the astrophysical and cosmological constraints upon its properties. However, before going into these specific topics, let us first note the strength of the basic big bang framework.

While Hubble’s work in the 1920’s established an expanding universe, the establishment of modern physical cosmology begins with the predictions of Gamow and his colleagues.

The basic mathematical space-time framework for the model that is now known as the Big Bang dates to the mid 1920s and the work of Alexander Friedman. However, Chandrasekhar has assured me that the prevailing cosmological picture in the 1930s was that the Hubble expansion started from a quasi-static Lemaitre-like model where galaxies never were squeezed together to form a different phase of matter. In particular, the hot big bang model hinges on two key quantitative observational tests: (1) the microwave background, and (2) big bang nucleosynthesis (BBN) and the light element abundances. This paper will focus on the second of these since that is more directly connected to high energy physics. However, it is worth noting that just as the new COBE<sup>1</sup> results have given renewed confidence in the  $3K$  background argument, the LEP collider (along with the SLC) has given us renewed confidence in the BBN arguments. We will return to this point momentarily. Note also that the microwave background probes events at temperatures  $\sim 10^4 K$  and times of  $\sim 10^5$  years, whereas the light element abundances probe the Universe at temperatures  $\sim 10^{10} K$  and times of  $\sim 1$  sec. Thus, it is the nucleosynthesis results that played the most significant role in leading to the particle-cosmology merger that has taken place this past decade.

Since the popular press sometimes presents misleading headlines implying *doubts* about the big bang, it is important to note here that the real concerns referred to in these articles are really in regard to observations related to models of galaxy and structure formation. The basic hot big bang model itself is in fantastic shape with high accuracy confirmations from COBE and, as we will discuss, nucleosynthesis. However, there is admittedly no fully developed model for galaxy and structure formation that fits all of the observations. (But, of course, there is also no fully developed first principles model for star formation either.) That we might not really know exactly how to make galaxies and large-scale structure in no way casts doubt on the hot, dense early universe which we call the big bang. (We also have trouble predicting earthquakes and tornadoes, but that hasn’t meant that we question celestial mechanics or a round Earth.) We will return to the problems of galaxy and structure formation towards the end of these lectures.

Before going into the specific argument as to the relationship of BBN to neutrino counting, let us review the history of BBN. This will draw heavily on other recent conference proceedings.<sup>2</sup>

## History of Big Bang Nucleosynthesis

It should be noted that there is a symbiotic connection between BBN and the  $3K$  background dating back to Gamow and his associates, Alpher and Herman. The initial

BBN calculations of Gamow and his associates<sup>3</sup> assumed pure neutrons as an initial condition and thus were not particularly accurate, but their inaccuracies had little effect on the group's predictions for a background radiation.

Once Hayashi recognized in 1950 the role of neutron-proton equilibration, the framework for BBN calculations themselves has not varied significantly. The work of Alpher, Follin and Herman<sup>4</sup> and Taylor and Hoyle,<sup>5</sup> preceeding the discovery of the 3K background, and Peebles<sup>6</sup> and Wagoner, Fowler and Hoyle,<sup>7</sup> immediately following the discovery, and the more recent work of our group of collaborators<sup>8,9,10,11,12</sup> all do essentially the same basic calculation, the results of which are shown in Figure 1. As far as the calculation itself goes, solving the reaction network is relatively simple by the standards of explosive nucleosynthesis calculations in supernovae (c.f. the 1965 calculation of Truran *et al.*),<sup>13</sup> with the changes over the last 25 years being mainly in terms of more recent nuclear reaction rates as input, not as any great calculational insight (although the current Kawano/Walker code<sup>11,12</sup> is somewhat streamlined relative to the earlier Wagoner code<sup>7</sup>). With the possible exception of  ${}^7\text{Li}$  yields, the reaction rate changes over the past 25 years have not had any major affect.<sup>9,11,12,13</sup> The one key improved input is a better neutron lifetime determination, a point to which we will also return shortly.

With the exception of the effects of elementary particle assumptions to which we will also return, the real excitement for BBN over the last 25 years has not really been in redoing the basic calculation. Instead, the true action is focused on understanding the evolution of the light element abundances and using that information to make powerful conclusions. In particular, in the 1960's, the main focus was on  ${}^4\text{He}$  which is very insensitive to the baryon density. The agreement between BBN predictions and observations helped support the basic big bang model but gave no significant information at that time with regard to density. In fact, in the mid-1960's, the other light isotopes (which are, in principle, capable of giving density information) were generally assumed to have been made during the T-Tauri phase of stellar evolution,<sup>15</sup> and so, were not then taken to have cosmological significance. It was during the 1970's that BBN fully developed as a tool for probing the universe. This possibility was in part stimulated by Ryter *et al.*<sup>16</sup> who showed that the T-Tauri mechanism for light element synthesis failed. Furthermore,  ${}^2\text{H}$  abundance determinations<sup>17,18</sup> improved significantly with solar wind measurements and the interstellar work from the Copernicus satellite. Reeves, Audouze, Fowler and Schramm<sup>19</sup> argued for cosmological  ${}^2\text{H}$  and were able to place a constraint on the baryon density excluding a universe closed with baryons. Subsequently, the  ${}^2\text{H}$  arguments were cemented when Epstein, Lattimer and Schramm<sup>20</sup> proved that no realistic astrophysical process other than the big bang could produce significant  ${}^2\text{H}$ . It was also interesting that the baryon density implied by BBN was in good agreement with the density implied by the dark galactic halos.<sup>21</sup>

By the late 1970's, a complimentary argument to  ${}^2\text{H}$  had also developed using  ${}^3\text{He}$ . In particular, it was argued<sup>22</sup> that, unlike  ${}^2\text{H}$ ,  ${}^3\text{He}$  was made in stars; thus, its abundance would increase with time. Since  ${}^3\text{He}$  like  ${}^2\text{H}$  monotonically decreased with cosmological baryon density, this argument could be used to place a lower limit on the baryon density<sup>23</sup> using  ${}^3\text{He}$  measurements from solar wind<sup>17</sup> or interstellar determinations.<sup>24</sup> Since the bulk of the  ${}^2\text{H}$  was converted in stars to  ${}^3\text{He}$ , the constraint was shown to be quite restrictive.<sup>9</sup> Support for this point<sup>25</sup> also comes from the observation of  ${}^3\text{He}$  in horizontal branch stars which, as processed stars still having  ${}^3\text{He}$  on their surface, indicates the survivability of

# BIG BANG NUCLEOSYNTHESIS

Walker, Steigman, Olive,  
and Schramm (1990)

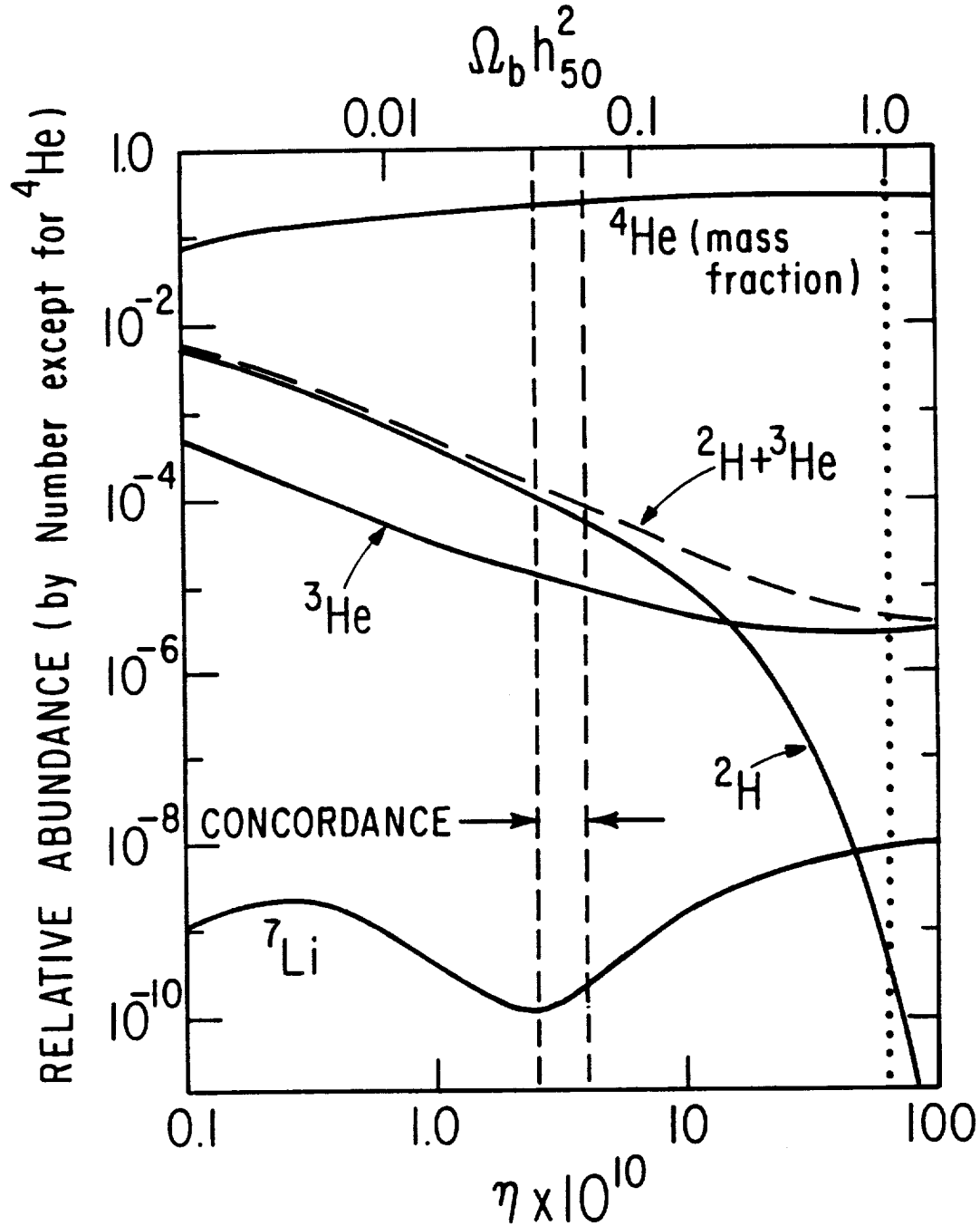


Figure 1. BBN abundances versus the baryon to photon ratio,  $\eta$ , or equivalently the fraction of the critical density,  $\Omega_b$ .

$^3\text{He}$ .

It was interesting that the lower boundary from  $^3\text{He}$  and the upper boundary from  $^2\text{H}$  yielded the requirement that  $^7\text{Li}$  be near its minimum of  $^7\text{Li}/\text{H} \sim 10^{-10}$ , which was verified by the Pop II  $\text{Li}$  measurements of Spite and Spite,<sup>26</sup> hence, yielding the situation emphasized by Yang *et al.*<sup>9</sup> that the light element abundances are consistent over nine orders of magnitude with BBN, but only if the cosmological baryon density is constrained to be around 6% of the critical value. It is worth noting that  $^7\text{Li}$  alone gives both an upper and a lower limit to  $\Omega_b$ . However, while its derived upper limit is more than competitive with the  $^2\text{H}$  limit, the  $^7\text{Li}$  lower limit is not nearly as restrictive as the  $^2\text{H} + ^3\text{He}$  limit. Claims that big bang nucleosynthesis can yield  $\Omega_b$  lower than 0.01 must necessarily neglect the  $^3\text{He} + ^2\text{H}$  limit.

The other development of the 70's for BBN was the explicit calculation of Steigman, Schramm and Gunn,<sup>27</sup> showing that the number of neutrino generations,  $N_\nu$ , had to be small to avoid overproduction of  $^4\text{He}$ . This will subsequently be referred to as the SSG limit. (Earlier work had noted a dependency of the  $^4\text{He}$  abundance on assumptions about the fraction of the cosmological stress-energy in exotic particles,<sup>28,5</sup> but had not actually made an explicit calculation probing the quantity of interest to particle physicists,  $N_\nu$ .) To put this in perspective, one should remember that the mid-1970's also saw the discovery of charm, bottom and tau, so that it almost seemed as if each new detector produced new particle discoveries, and yet, cosmology was arguing against this "conventional" wisdom. Over the years the SSG limit on  $N_\nu$  improved with  $^4\text{He}$  abundance measurements, neutron lifetime measurements and with limits on the lower bound to the baryon density; hovering at  $N_\nu \lesssim 4$  for most of the 1980's and dropping to slightly lower than 4<sup>29,30,10</sup> just before LEP and SLC turned on.

## Big Bang Nucleosynthesis: $\Omega_b$ and $N_\nu$

The power of big bang nucleosynthesis comes from the fact that essentially all of the physics input is well determined in the terrestrial laboratory. The appropriate temperature regimes, 0.1 to 1 MeV, are well explored in nuclear physics labs. Thus, what nuclei do under such conditions is not a matter of guesswork, but is precisely known. In fact, it is known for these temperatures far better than it is for the centers of stars like our sun. The center of the sun is only a little over 1 keV, thus, below the energy where nuclear reaction rates yield significant results in laboratory experiments, and only the long times and higher densities available in stars enable anything to take place.

To calculate what happens in the big bang, all one has to do is follow what a gas of baryons with density  $\rho_b$  does as the universe expands and cools. As far as nuclear reactions are concerned, the only relevant region is from a little above 1 MeV ( $\sim 10^{10} \text{K}$ ) down to a little below 100 keV ( $\sim 10^9 \text{K}$ ). At higher temperatures, no complex nuclei other than free single neutrons and protons can exist, and the ratio of neutrons to protons,  $n/p$ , is just determined by

$$n/p = e^{-Q/T},$$

where  $Q = (m_n - m_p)c^2 \sim 1.3 \text{MeV}$ .

Equilibrium applies because the weak interaction rates are much faster than the expansion of the universe at temperatures much above  $10^{10} \text{K}$ . At temperatures much below  $10^9 \text{K}$ ,

the electrostatic repulsion of nuclei prevents nuclear reactions from proceeding as fast as the cosmological expansion separates the particles.

Because of the equilibrium existing for temperatures much above  $10^{10}K$ , we don't have to worry about what went on in the universe at higher temperatures. Thus, we can start our calculation at  $10MeV$  and not worry about speculative physics like the theory of everything (T.O.E.), or grand unifying theories (GUTs), as long as a gas of neutrons and protons exists in thermal equilibrium by the time the universe has cooled to  $\sim 10MeV$ .

After the weak interaction drops out of equilibrium, a little above  $10^{10}K$ , the ratio of neutrons to protons changes more slowly due to free neutrons decaying to protons, and similar transformations of neutrons to protons via interactions with the ambient leptons. By the time the universe reaches  $10^9K$  ( $0.1MeV$ ), the ratio is slightly below  $1/7$ . For temperatures above  $10^9K$ , no significant abundance of complex nuclei can exist due to the continued existence of gammas with greater than  $MeV$  energies. Note that the high photon to baryon ratio in the universe ( $\sim 10^{10}$ ) enables significant population of the  $MeV$  high energy Boltzman tail until  $T \lesssim 0.1 MeV$ .

Once the temperature drops to about  $10^9K$ , nuclei can exist in statistical equilibrium through reactions such as  $n + p \leftrightarrow {}^2H + \gamma$  and  ${}^2H + p \leftrightarrow {}^3He + \gamma$  and  ${}^2H + n \leftrightarrow {}^3H + \gamma$ , which in turn react to yield  ${}^4He$ . Since  ${}^4He$  is the most tightly bound nucleus in the region, the flow of reactions converts almost all the neutrons that exist at  $10^9K$  into  ${}^4He$ . The flow essentially stops there because there are no stable nuclei at either mass-5 or mass-8. Since the baryon density at big bang nucleosynthesis is relatively low (much less than  $1g/cm^3$ ) and the time-scale short ( $t \lesssim 10^2sec$ ), only reactions involving two-particle collisions occur. It can be seen that combining the most abundant nuclei, protons and  ${}^4He$  via two body interactions always leads to unstable mass-5. Even when one combines  ${}^4He$  with rarer nuclei like  ${}^3H$  or  ${}^3He$ , we still get only to mass-7, which, when hit by a proton, the most abundant nucleus around, yields mass-8. (A loophole around the mass-8 gap can be found if  $n/p > 1$ , so that excess neutrons exist, but for the standard case  $n/p < 1$ ). Eventually,  ${}^3H$  radioactively decays to  ${}^3He$ , and any mass-7 made radioactively decays to  ${}^7Li$ . Thus, big bang nucleosynthesis makes  ${}^4He$  with traces of  ${}^2H$ ,  ${}^3He$ , and  ${}^7Li$ . (Also, all the protons left over that did not capture neutrons remain as hydrogen.) For standard homogeneous BBN, all other chemical elements are made later in stars and in related processes. (Stars jump the mass-5 and -8 instability by having gravity compress the matter to sufficient densities and have much longer times available so that three-body collisions can occur.) With the possible exception of  ${}^7Li$ ,<sup>9,10,11,12,14</sup> the results are rather insensitive to the detailed nuclear reaction rates. This insensitivity was discussed in reference 9 and most recently using a Monte Carlo study by Krauss and Romanelli.<sup>14</sup> An  $n/p$  ratio of  $\sim 1/7$  yields a  ${}^4He$  primordial mass fraction,

$$Y_p = \frac{2n/p}{n/p + 1} \approx \frac{1}{4}.$$

The only parameter we can easily vary in such calculations is the density that corresponds to a given temperature. From the thermodynamics of an expanding universe we know that  $\rho_b \propto T^3$ ; thus, we can relate the baryon density at  $10^{10}K$  to the baryon density today, when the temperature is about  $3K$ . The problem is that we don't know today's  $\rho_b$ , so the calculation is carried out for a range in  $\rho_b$ . Another aspect of the density is that the cosmological expansion rate depends on the total mass-energy density associated with a

given temperature. For cosmological temperatures much above  $10^4 K$ , the energy density of radiation exceeds the mass-energy density of the baryon gas. Thus, during big bang nucleosynthesis, we need the radiation density as well as the baryon density. The baryon density determines the density of the nuclei and thus their interaction rates, and the radiation density controls the expansion rate of the universe at those times. The density of radiation is just proportional to the number of types of radiation. Thus, the density of radiation is not a free parameter if we know how many types of relativistic particles exist when big bang nucleosynthesis occurred.

Assuming that the allowed relativistic particles at  $1 MeV$  are photons,  $e, \mu$ , and  $\tau$  neutrinos (and their antiparticles) and electrons (and positrons), Figure 1 shows the BBN yields for a range in present  $\rho_b$ , going from less than that observed in galaxies to greater than that allowed by the observed large-scale dynamics of the universe. The  ${}^4He$  yield is almost independent of the baryon density, with a very slight rise in the density due to the ability of nuclei to hold together at slightly higher temperatures and at higher densities, thus enabling nucleosynthesis to start slightly earlier, when the baryon to photon ratio is higher. No matter what assumptions one makes about the baryon density, it is clear that  ${}^4He$  is predicted by big bang nucleosynthesis to be around 1/4 of the mass of the universe.

## The SSG Limit - Cosmological Neutrino Counting

Let us now look at the connection to  $N_\nu$ . Remember that the yield of  ${}^4He$  is very sensitive to the  $n/p$  ratio. The more types of relativistic particles, the greater the energy density at a given temperature, and thus, a faster cosmological expansion. A faster expansion yields the weak-interaction rates being exceeded by the cosmological expansion rate at an earlier, higher temperature; thus, the weak interaction drops out of equilibrium sooner, yielding a higher  $n/p$  ratio. It also yields less time between dropping out of equilibrium and nucleosynthesis at  $10^9 K$ , which gives less time for neutrons to change into protons, thus also increasing the  $n/p$  ratio. A higher  $n/p$  ratio yields more  ${}^4He$ . As we will see in the next section, quark-hadron induced variations<sup>31</sup> in the standard model also yield higher  ${}^4He$  for higher values of  $\Omega_b$ . Thus, such variants still support the constraint on the number of relativistic species.<sup>32</sup>

In the standard calculation we allowed for photons, electrons, and the three known neutrino species (and their antiparticles). However, following SSG and doing the calculation (see Figure 2) for additional species of neutrinos, we can see when  ${}^4He$  yields exceed observational limits while still yielding a density consistent with the  $\rho_b$  bounds from  ${}^2H$ ,  ${}^3He$ , and now  ${}^7Li$ . (The new  ${}^7Li$  value gives approximately the same constraint on  $\rho_b$  as the others, thus strengthening the conclusion.) The bound on  ${}^4He$  comes from observations of helium in many different objects in the universe. However, since  ${}^4He$  is not only produced in the big bang but in stars as well, it is important to estimate what part of the helium in some astronomical object is primordial—from the big bang—and what part is due to stellar production after the big bang. The pioneering work of the Peimberts<sup>33</sup> showing that  ${}^4He$  varies with oxygen has now been supplemented by examination of how  ${}^4He$  varies with nitrogen and carbon. The observations have also been systematically re-examined by Pagel.<sup>34</sup> The conclusions of Pagel,<sup>34</sup> Steigman *et al.*<sup>35</sup> and Walker *et al.*<sup>11</sup> all agree that the  ${}^4He$  mass fraction,  $Y_p$ , extrapolated to zero heavy elements, whether using  $N$ ,  $O$ , or  $C$ , is  $Y_p \sim 0.23$  with an upper bound allowing for possible systematics of 0.24.

The other major uncertainty in the  ${}^4He$  production used to be the neutron lifetime.

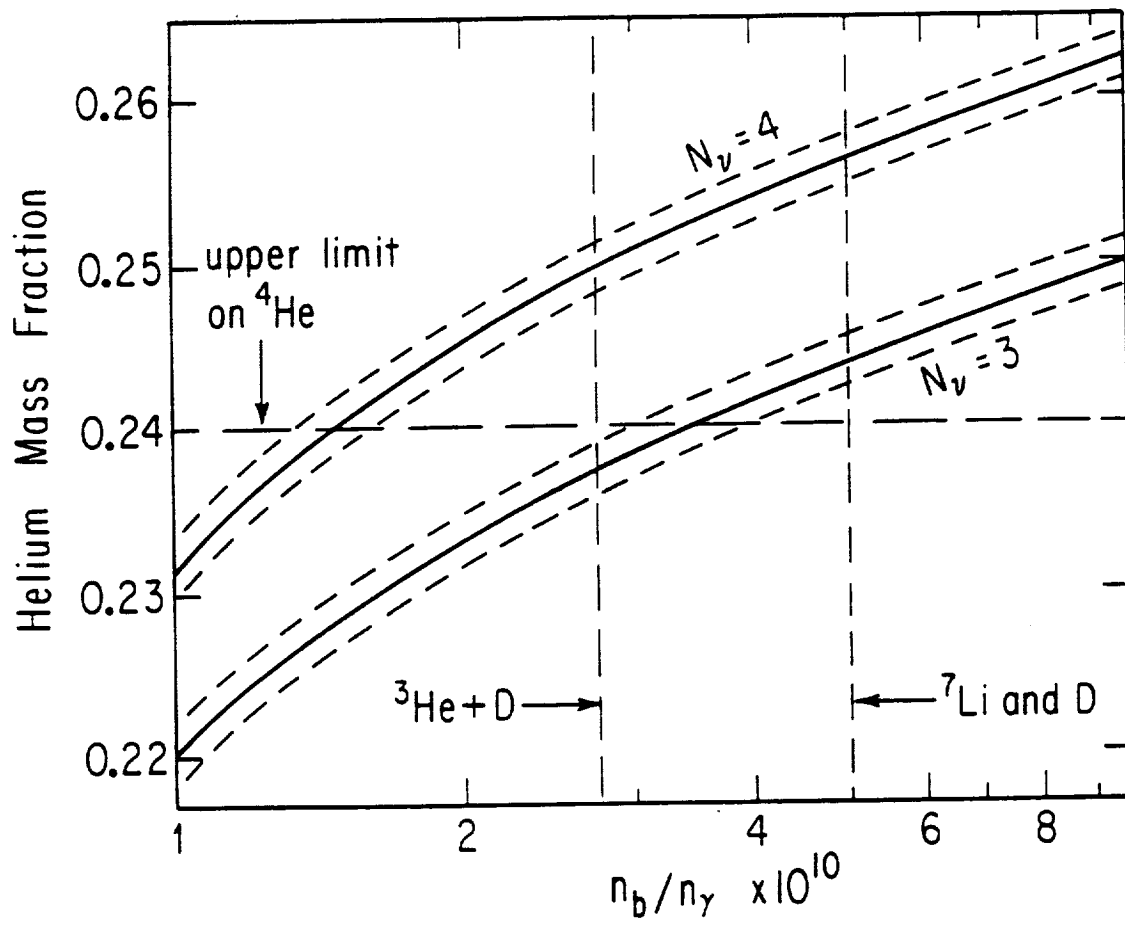


Figure 2. The SSG argument with recent parameter constraints showing the BBN helium mass fraction versus  $\eta$  for  $N_\nu = 3$  and 4. Note that 4 is excluded.



However, the new world average of  $\tau_n = 890 \pm 4s$  ( $\tau_{1/2} = 10.3 \text{ min}$ ) is dominated by the dramatic results of Mampe *et al.*<sup>36</sup> using a neutron bottle. This new result is quite consistent with a new counting measurement of Byrne *et al.*<sup>37</sup> and within the errors of the previous world average of  $896 \pm 10s$  and is also consistent with the precise  $C_A/C_V$  measurements from PERKEO<sup>38</sup> and others. Thus, the old ranges of  $10.4 \pm 0.2 \text{ min}$ , used for the half-life in calculations,<sup>39,9</sup> seem to have converged towards the lower side. The convergence means that, instead of the previous broad bands for each neutrino flavor, we obtain relatively narrow bands (see Figure 2). Note that  $N_\nu = 4$  is excluded. In fact, the SSG limit is now  $N_\nu < 3.4$ .<sup>10,11</sup>

The recent verification of this cosmological standard model prediction by LEP,  $N_\nu = 2.98 \pm 0.06$ , from the average of ALEPH, DELPHI, L3 and OPAL<sup>40</sup> collaborations as well as the SLC<sup>40</sup> results, thus experimentally confirms our confidence in BBN. (However, we should also remember that LEP and cosmology are sensitive to different things.<sup>41</sup> Cosmology counts all relativistic degrees of freedom for  $m_x \lesssim 10 \text{ MeV}$ , with LEP and SLC counting particles coupling to the  $Z^0$  with  $m_x \lesssim 45 \text{ GeV}$ .)

While  $\nu_e$  and  $\nu_\mu$  are obviously counted equally in both situations, a curious loophole exists for  $\nu_\tau$  since the current experimental limit  $m_{\nu_\tau} < 35 \text{ MeV}$  could allow it not to contribute as a full neutrino in the cosmology argument.<sup>42</sup> Proposed experiments which push the  $m_{\nu_\tau}$  limit down to less than a few  $\text{MeV}$  should eliminate this loophole. It might also be noted that if we assume  $m_{\nu_\tau}$  is light so that cosmologically  $N_\nu = 3$ , we can turn the argument around and use LEP to predict the primordial helium abundance ( $\sim 24\%$ ), or even use limits on  ${}^4\text{He}$  to give an upper limit on  $\Omega_b$  (also  $\lesssim 0.10$ ). Thus, LEP strengthens the argument that we need non-baryonic dark matter if  $\Omega = 1$ . In fact, note also that with  $N_\nu = 3$ , if  $Y_p$  is ever proven to be less than  $\sim 0.235$ , standard BBN is in difficulty. Similar difficulties occur if  $Li/H$  is ever found below  $\sim 10^{-10}$ . In other words, BBN is a falsifiable theory.

## Alternative Proposals

As noted above, BBN yields all agree with observations using only one freely adjustable parameter,  $\rho_b$ . Thus, BBN can make strong statements regarding  $\rho_b$  if the observed light element abundances cannot be fit with any alternative theory. Before exploring the implications for  $\rho_b$ , let us examine alternative proposals which have arisen to try to escape the power of the homogeneous BBN conclusions.

The two alternatives that have recently received interest are:

- (1) Decaying particles;<sup>43</sup> and
- (2) Quark-hadron transition inspired inhomogeneities.<sup>31</sup>

The first of these notes that if a species of massive particle ( $M \gtrsim \text{few GeV}$ ) were to decay after traditional BBN, it could redo nucleosynthesis. While previous decaying particle proposals had been made, the new idea<sup>43</sup> emphasizes the importance of the resulting hadron cascade which, they argue, will dominate the yields. While interesting results are obtained, problems with detailed abundance determinations do result. In particular, this class of models seems to predict inevitably that  ${}^6\text{Li}/{}^7\text{Li} \gg 1$ , whereas observations show  ${}^7\text{Li}/{}^6\text{Li} \gtrsim 10$ . While at first this might seem fatal, it is almost avoidable by noting that  ${}^6\text{Li}$  is much more fragile than  ${}^7\text{Li}$ ; thus, it is easy to deplete  ${}^6\text{Li}$  and obtain the observed ratios. However, Brown and Schramm<sup>44</sup> have pointed out that for high surface temperature Pop II stars, the convective zones do not go deep enough to destroy any primordial

${}^6\text{Li}$ . Pilachowski *et al.*<sup>45</sup> have now looked at those specific stars and indeed find no  ${}^6\text{Li}$ , again seeing  ${}^7\text{Li}/{}^6\text{Li} > 10$ . Therefore, unless the Brown and Schramm convection argument can be surmounted,  ${}^6\text{Li}$  seems to constrain this model seriously. Steigman, Audouze and others have noted additional problems with this model for  ${}^3\text{He}$  and  ${}^2\text{H}$  ratios.

Let us now look at the quark-hadron inspired inhomogeneity models.<sup>31</sup> While inhomogeneity models had been looked at previously (c.f. reference 9) and were found to make little difference, the quark-hadron inspired models had the added ingredient of variations in  $n/p$  ratios.

The initial claim by Applegate *et al.*, followed by a similar argument from Alcock *et al.* that  $\Omega_b \sim 1$  might be possible, created tremendous interest. Their argument was that if the quark-hadron transition was a first-order phase transition (as some preliminary lattice gauge calculations implied), then it was possible that large inhomogeneities could develop at  $T \gtrsim 100\text{MeV}$ . The preferential diffusion of neutrons versus protons out of the high density regions could lead to big bang nucleosynthesis occurring under conditions with both density inhomogeneities and variable neutron/proton ratios. In the first round of calculations, it was claimed that such conditions might allow  $\Omega_b \sim 1$ , while fitting the observed primordial abundances of  ${}^4\text{He}$ ,  ${}^2\text{H}$ ,  ${}^3\text{He}$  with an overproduction of  ${}^7\text{Li}$ . Since  ${}^7\text{Li}$  is the most recent of the cosmological abundance constraints and has a different observed abundance in Pop I stars versus the traditionally more primitive Pop II stars,<sup>26</sup> some argued that perhaps some special depletion process might be going on to reduce the excess  ${}^7\text{Li}$ . Reeves and Audouze each argued against such processes and tried to turn the argument around and use lithium abundances to constrain the quark-hadron transition.

At first it appeared that if the lithium constraint could be surmounted, then the constraints of standard big bang nucleosynthesis might disintegrate. (Although Audouze, Reeves and Schramm emphasized that the number of parameters needed to fit the light elements was somewhat larger for these non-standard models, nonetheless, a non-trivial loophole appeared to be forming.) To further stimulate the flow through the loophole, Mullaney and Fowler showed that, in addition to looking at the diffusion of neutrons out of high density regions, one must also look at the subsequent effect of excess neutrons diffusing back into the high density regions as the nucleosynthesis goes to completion in the low density regions. (The initial calculations treated the two regions separately.) Mullaney and Fowler argued that for certain phase transition parameter values (e.g. nucleation site separations  $\sim 10m$  at the time of the transition), this back diffusion could destroy much of the excess lithium. Recent work by Banerjee and Chitre (private communication) suggests that more accurate treatment of the diffusion calculation could reduce the interesting separation distance by several orders of magnitude.

However, Kurki-Suonio, Matzner, Olive and Schramm,<sup>32</sup> the Tokyo group,<sup>46</sup> and the Livermore group<sup>47</sup> have recently argued that in their detailed diffusion models, the back diffusion not only effects  ${}^7\text{Li}$ , but also the other light nuclei as well. They find that for  $\Omega_b \sim 1$ ,  ${}^4\text{He}$  is also overproduced (although it does go to a minimum for similar parameter values as does the lithium). One can understand why these models might tend to overproduce  ${}^4\text{He}$  and  ${}^7\text{Li}$  by remembering that in standard homogeneous big bang nucleosynthesis, high baryon densities lead to excesses in these nuclei. As back diffusion evens out the effects of the initial fluctuation, the averaged result should approach the homogeneous value. Furthermore, it can be argued that any narrow range of parameters, such as those which yield relatively low lithium and helium, are unrealistic since in most realistic phase

transitions there are distributions of parameter values (distribution of nucleation sites, separations, density fluctuations, etc.). Therefore, narrow minima are washed out which would bring the  ${}^7\text{Li}$  and  ${}^4\text{He}$  values back up to their excessive levels for all parameter values with  $\Omega \sim 1$ . Furthermore, Adams and Freese<sup>48</sup> have argued that the boundary between the two phases may be fractal-like rather than smooth. The large surface area of a fractal-like boundary would allow more interaction between the regions and minimize exotic effects.

Figure 3 shows the results of Kurki-Suonio *et al.*<sup>32</sup> for varying spacing  $l$  with the constraints from the different light element abundances. Notice that the  $\text{Li}$  and even the  ${}^4\text{He}$  constraint do not allow  $\Omega_b \sim 1$ . (The  ${}^4\text{He}$  abundance constraint used in Kurki-Suonio *et al.* was a generous  $Y_p \lesssim 0.25$ ; for the preferred  $Y_p \lesssim 0.24$ , the  ${}^4\text{He}$  bound is about as tight as the Pop II  $\text{Li}$  constraint.) Note also that with the Pop II  ${}^7\text{Li}$  constraint, the results for  $\Omega_b$  are quite similar to the standard model with a slight excess in  $\Omega_b$  possible if  $l$  is tuned to  $\sim 10$ .

Furthermore, initially it looked like quark-hadron inspired models might enable leakage<sup>49</sup> beyond mass-7, thus enabling  ${}^9\text{Be}$ ,  ${}^{14}\text{N}$ , or maybe even r-process elements to become probes as whether or not the universe had such a transition (even if  $\Omega_b \sim 1$ ). However, Tarasawa and Sato<sup>46</sup> have shown that when full multizone calculations of the type used by Kurki-Suonio *et al.* are utilized, then no significant leakage occurs.

One possible signature that remains for a first order quark-hadron transition is a slightly larger allowed range for  $Y_p$  that is concordant with  $N_\nu = 3$  and with the other light element abundances. In particular, if  ${}^4\text{He}$  were ever shown to be definitively  $\lesssim 0.23$ , it might be evidence for such a quark-hadron induced behavior since the standard homogeneous case cannot accomodate such values. Of course, excessively low values for  $Y_p$  would still be unallowable.

One can conclude from the failure of the attempts to circumvent the standard BBN results that the results are amazingly robust. Even when many new free parameters are added, as in the quark-hadron case, the bottom line, when one requires concordance with the light element abundances, is essentially the same as the standard result. In other words,  $\Omega_b \sim 0.06$  (although with fine-tuning the upper bound might be relaxed a bit to  $\sim 0.2$  rather than 0.1).

One loophole which can yield variations in  $\Omega_b$  outside the above range is to allow for degenerate neutrinos. This possibility has been discussed by many authors over the years with the most recent being Olive *et al.*<sup>121</sup> (also see references therein). The basic point is that creating any significant deviations requires excess lepton number densities,  $Li$ , that are comparable to  $\eta_\gamma$ . However, most grand unified/SUSY theories require Lepton number excesses that are comparable to baryon number excess. Thus,  $Li \sim \eta_B \ll \eta_\gamma$ .

Furthermore, in order to fit the observed abundances, (word missing) requires that  $Le \neq L_{\mu,\tau}$ . Since again unified theories tend not to produce  $Le$  vs.  $L_{\mu,\tau}$  in the required ratios, but more like  $Le \sim L_\mu \sim L_\tau \sim \eta_b$ , this loophole appears rather unnatural and requires additional parameters that require artificial tuning.

## Limits on $\Omega_b$ and Dark Matter Requirements

The narrow range in baryon density for which concordance occurs is very interesting. Let us convert it into units of the critical cosmological density for the allowed range of Hubble expansion rates. For the big bang nucleosynthesis constraints,<sup>9,10,11,12,29,30</sup> the

# "QUARK-HADRON" NUCLEOSYNTHESIS

Kurki-Suonio, Matzner, Olive and Schramm (1990)

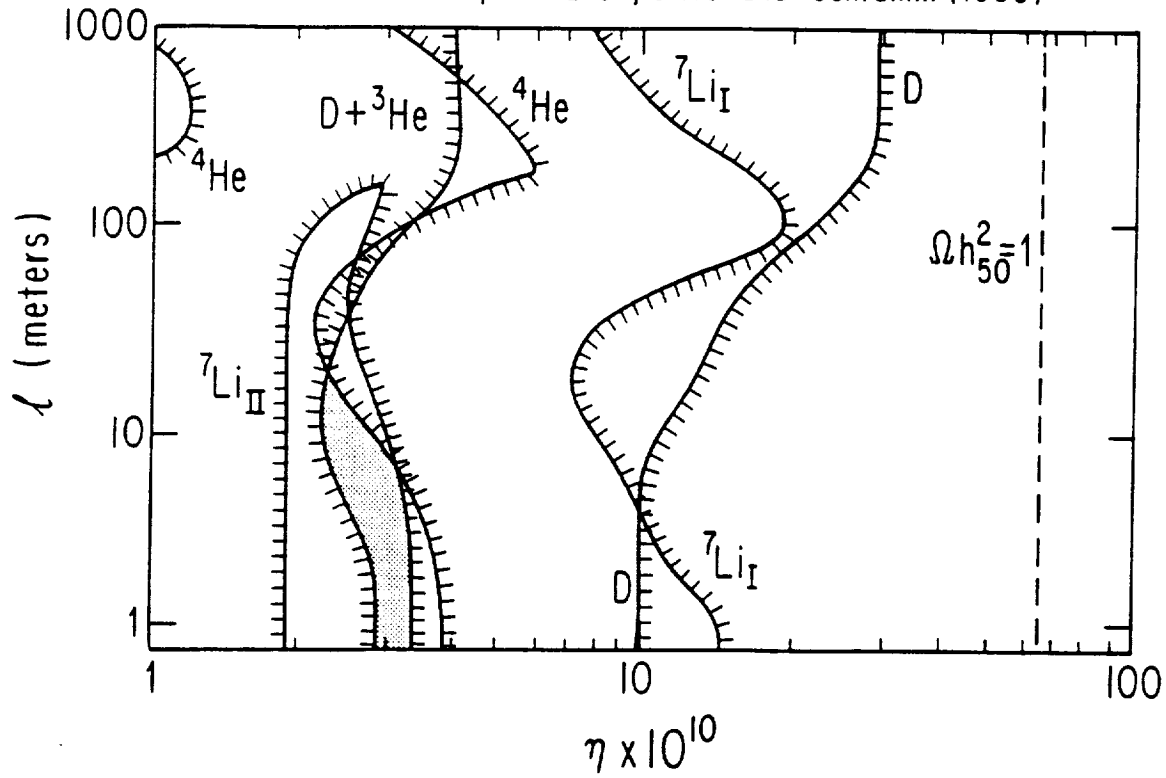


Figure 3. This shows the constraints on  $\eta$  of the various observed abundances in a first-order quark-hadron phase transition with nucleation sites separated by a distance  $l$  with density contrast  $R \lesssim 10^3$ . The Pop II lithium abundance used here is from the compilation of data given by Walker *et al.*<sup>11</sup> and is slightly more restrictive on  $\eta$  than that used in Figure 2 or used in the original Kurki-Suonio *et al.*<sup>32</sup> calculation from which this figure is derived. It should be noted that work by the Tokyo<sup>46</sup> group and by the Livermore group<sup>47</sup> confirms the conclusions on restricting  $\Omega_b$  to values similar to the standard result even when  $R \rightarrow \infty$ .

dimensionless baryon density  $\Omega_b$ , that fraction of the critical density that is in baryons, is less than 0.11 and greater than 0.02 for  $0.04 \lesssim h_0 \lesssim 0.7$ , where  $h_0$  is the Hubble constant in units of  $100 \text{ km/sec/Mpc}$ . The lower bound on  $h_0$  comes from direct observational limits and the upper bound from age of the universe constraints.<sup>49</sup> Note that the constraint on  $\Omega_b$  means that the universe *cannot be closed with baryonic matter*. If the universe is truly at its critical density, then nonbaryonic matter is required. This argument has led to one of the major areas of research at the particle-cosmology interface, namely, the search for non-baryonic dark matter.

Another important conclusion regarding the allowed range in baryon density is that it is in very good agreement with the density implied from the dynamics of galaxies, *including their dark halos*. An early version of this argument, using only deuterium, was described over fifteen years ago.<sup>21</sup> As time has gone on, the argument has strengthened, and the fact remains that galaxy dynamics and nucleosynthesis agree at about 6% of the critical density. Thus, if the universe is indeed at its critical density, as many of us believe, it requires most matter not to be associated with galaxies and their halos, as well as to be nonbaryonic.

Let us put the nucleosynthetic arguments in context. The arguments requiring some sort of dark matter fall into two separate and quite distinct areas. First are the arguments using Newtonian mechanics applied to various astronomical systems that show that there is more matter present than the amount that is shining. These arguments are summarized in Figure 4. It should be noted that these arguments reliably demonstrate that galactic halos seem to have a mass  $\sim 10$  times the visible mass.

Note, however, that big bang nucleosynthesis requires that the bulk of the baryons in the universe be dark since  $\Omega_{vis} \ll \Omega_b$ . Thus, the dark halos could in principle be baryonic.<sup>21</sup> Recently, arguments on very large scales<sup>51</sup> (bigger than clusters of galaxies) hint that  $\Omega$  on those scales is indeed greater than  $\Omega_b$ , thus forcing us to need non-baryonic matter. This is the first observational support for an  $\Omega$  bigger than what can be accommodated with  $\Omega_b$ .

Of course, it has been long anticipated, since the only long-lived natural value for  $\Omega$  is unity, that inflation<sup>52</sup> or something like it provided the early universe with the mechanism to achieve that value and thereby solve the flatness and smoothness problems.

Some baryonic dark matter must exist, since from the  $^2\text{H} + ^3\text{He}$  argument we know that the lower bound from big bang nucleosynthesis is greater than the upper limits on the amount of visible matter in the universe. However, we do not know what form this baryonic dark matter is in. It could be either in condensed objects in the halo, such as brown dwarfs and jupiters (objects with  $\lesssim 0.08 M_\odot$  so they are not bright shining stars), or in black holes (which at the time of nucleosynthesis would have been baryons). Or, if the baryonic dark matter is not in the halo, it could be in hot intergalactic gas, hot enough not to show absorption lines in the Gunn-Peterson test, but not so hot as to be seen in the x-rays. Evidence for some hot gas is found in clusters of galaxies. However, the amount of gas in clusters would not be enough to make up the entire missing baryonic matter. Another possible hiding place for the dark baryons would be failed galaxies, large clumps of baryons that condense gravitationally but did not produce stars. Such clumps are predicted in galaxy formation scenarios that include large amounts of biasing where only some fraction of the clumps shine.

Hegyi and Olive<sup>53</sup> have argued that dark baryonic halos are unlikely. However, they do

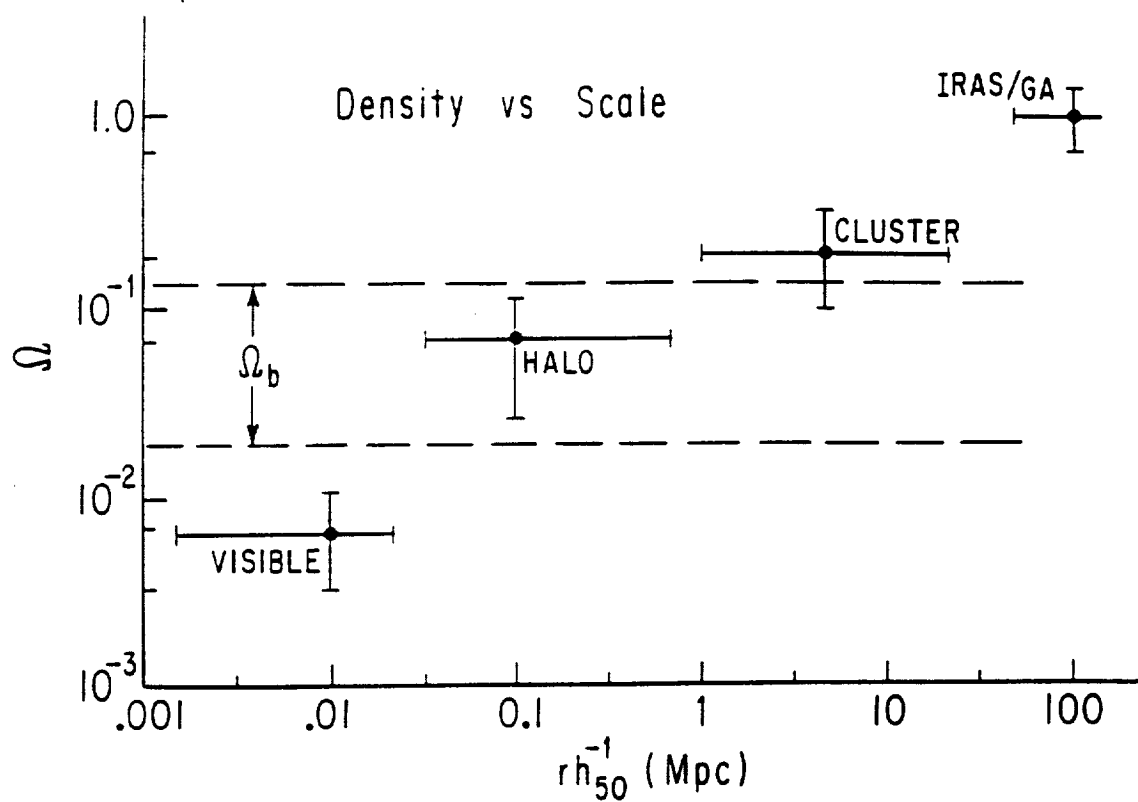


Figure 4. The inferred density in units of the critical density as a function of the scale on which it is "measured." Note the increase in  $\Omega$  towards unity as larger scales are probed. Note also that  $\Omega_b$  agrees with densities on the scale of galactic halos and is greater than the amount of visible matter.

Table 1

MATTER

Baryonic  $\Omega_b \sim 0.06$

VISIBLE  $\Omega_{vis} \lesssim 0.01$

DARK

Halo

Jupiters

Brown Dwarfs

Stellar Black Holes

Intergalactic

Hot gas at  $T \sim 10^5 K$

Stillborn Galaxies

Non Baryonic  $\Omega_{nb} \sim 0.94$

HOT

$m_{\nu_\tau} \sim 25 eV$

COLD

Wimps/Inos  $\sim 100 Gev$

Axions  $\sim 10^{-5} eV$

Planetary Mass Black Holes

allow for the loopholes mentioned above of low mass objects or of massive black holes. It is worth noting, as Schramm<sup>2</sup> points out, that these loopholes are not that unlikely. Furthermore, recent observational evidence<sup>54</sup> seems to show that disk formation is relatively late, occurring at red shifts  $z \lesssim 1$ . Thus, the first several billion years of a galaxy's life may have been spent prior to the formation of the disk. In fact, if the first large objects to form are less than galactic mass, as many scenarios imply, then mergers are necessary for eventual galaxy size objects. Mergers stimulate star formation while putting early objects into halos rather than disks. Mathews and Schramm<sup>55</sup> have recently developed a galactic evolution model which does just that and gives a reasonable scenario for chemical evolution. (This scenario also provides a natural explanation for the number-versus-redshift relation of low luminosity galaxies found by Cowie.<sup>56</sup> Thus, while making halos out of exotic material may be more exciting, it is certainly not impossible for the halos to be in the form of dark baryons. One application of William of Ockham's famous razor would be to have us not invoke exotic matter until we are forced to do so.

Non-baryonic matter can be divided following Bond and Szalay<sup>57</sup> into two major categories for cosmological purposes: hot dark matter (HDM) and cold dark matter (CDM). Hot dark matter is matter that is relativistic until just before the epoch of galaxy formation, the best example being low mass neutrinos with  $m_\nu \sim 25\text{eV}$ . (Remember  $\Omega_\nu \sim \frac{m_\nu(\text{eV})}{100h_0^2}$ ).

Cold dark matter is matter that is moving slowly at the epoch of galaxy formation. Because it is moving slowly, it can clump on very small scales, whereas HDM tends to have more difficulty in being confined on small scales. Examples of CDM could be massive neutrino-like particles with masses,  $M_x$ , greater than several  $\text{GeV}$  or the lightest supersymmetric particle which is presumed to be stable and might also have masses of several  $\text{GeV}$ . Following Michael Turner, all such weakly interacting massive particles are called "WIMPS." Axions, while very light, would also be moving very slowly<sup>58</sup> and, thus, would clump on small scales. Or, one could also go to non-elementary particle candidates, such as planetary mass blackholes or quark nuggets of strange quark matter, possibly produced at the quark-hadron transition.<sup>59</sup> Another possibility would be any sort of massive topological remnant left over from some early phase transition. Table 1 summarizes the matter options. Note that CDM would clump in halos, thus requiring the dark baryonic matter to be out between galaxies, whereas HDM would allow baryonic halos.

When thinking about dark matter candidates, one should remember the basic work of Zeldovich,<sup>60</sup> resurrected by Lee and Weinberg<sup>61</sup> and others,<sup>62</sup> which showed that for a weakly interacting particle, one can obtain closure densities, either if the particle is very light,  $\sim 25\text{eV}$ , or if the particle is very massive,  $\sim 3\text{GeV}$ . This occurs because, if the particle is much lighter than the decoupling temperature, then its number density is the number density of photons (to within spin factors and small corrections), and so the mass density is in direct proportion to the particle mass, since the number density is fixed. However, if the mass of the particle is much greater than the decoupling temperature, then annihilations will deplete the particle number since, as the temperature of the expanding universe drops below the rest mass of the particle, Boltzmann suppression prohibits production while the number is depleted via annihilations until the annihilation reaction freezes out. For normal weakly interacting particles, decoupling occurs at a temperature of  $\sim 1\text{MeV}$ , so higher mass particles are depleted. It should also be noted that the curve of density versus particle mass turns over again (see Figure 5) once the mass of the WIMP exceeds the mass of the coupling boson<sup>63,64,65</sup> so that the annihilation cross section varies



# Zeldovich - Lee - Weinberg - etc Argument

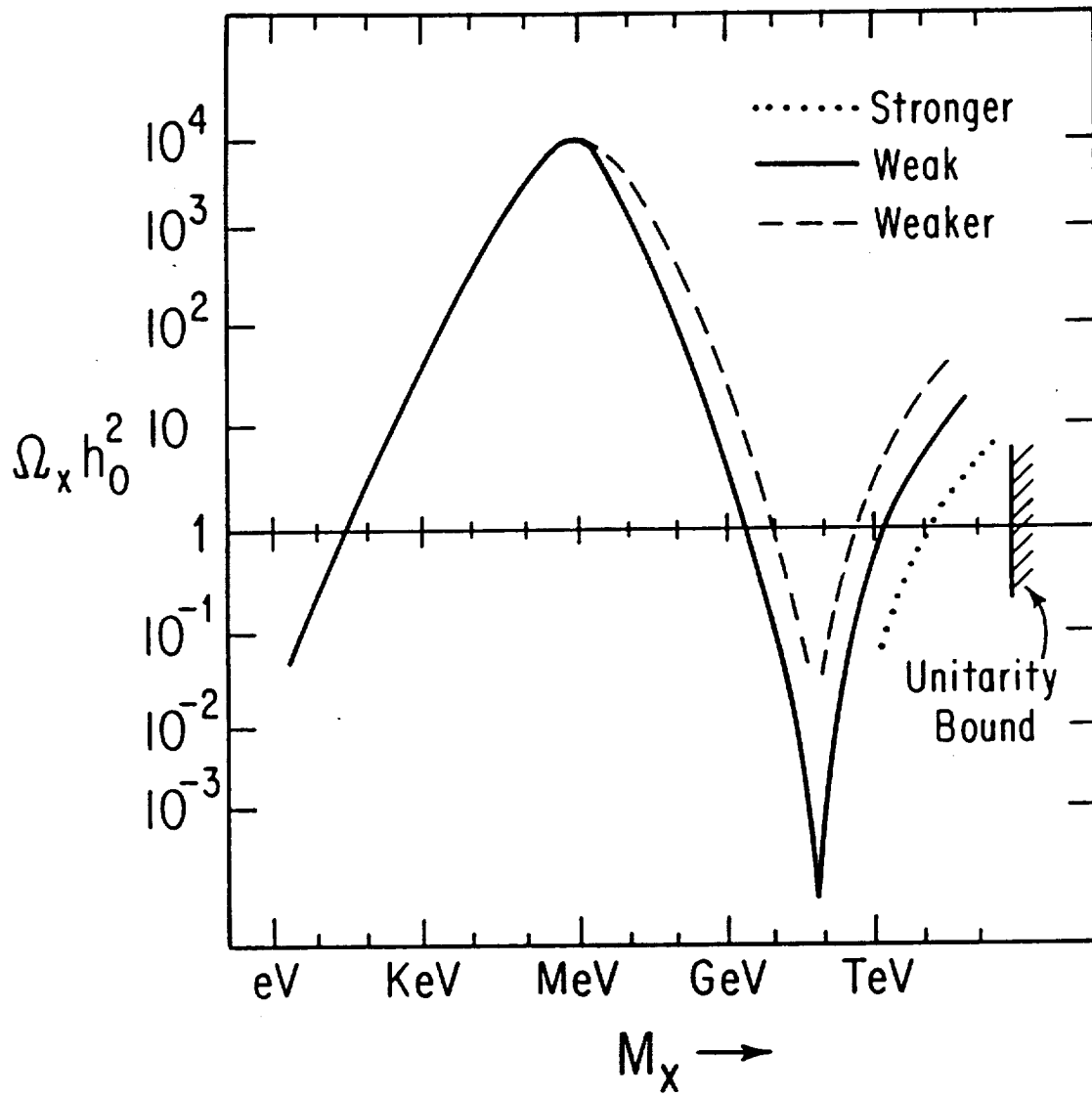


Figure 5.  $\Omega_\chi h_0^2$  versus  $M_\chi$  for weakly interacting particles showing three crossings of  $\Omega h_0^2 = 1$ . Note also how the curve shifts at high  $M_\chi$  for interactions weaker or stronger than normal weak interaction (where normal weak is that of neutrino coupling through  $Z^0$ ). Extreme strong couplings reach a unitarity limit at  $M_\chi \sim 340 \text{ TeV}$ .

as  $\frac{1}{M_x^2}$ , independent of the mass of the coupling boson. In this latter case,  $\Omega = 1$  can be obtained for  $M_x \sim 1\text{TeV} \sim (3K \times M_{\text{Planck}})^{1/2}$ , where  $3K$  and  $M_{\text{Planck}}$  are the only energy scales left in the calculation (see Figure 5). A loophole to this argument occurs if there is a matter-antimatter asymmetry as in the case of baryons. However, such particles would have to be Dirac particles and we will see that they are still severely constrained.

A few years ago the preferred candidate particle was probably a few  $\text{GeV}$  mass WIMP. However, LEP's lack of discovery of any new particle coupling to the  $Z^0$  with  $M_x \lesssim 45\text{GeV}$ , coupled with underground experiments,<sup>66</sup> clearly eliminates that candidate.<sup>67,68</sup> Constraints for particles not fully coupled to the  $Z^0$  were discussed by Ellis, Nanopoulos, Roskowski and Schramm<sup>68</sup> and are updated and presented in Figures 6a and 6b. (The inclusion of the Kamiokande II results as well as the newer LEP limits yields an important update over the results of Ellis *et al.*<sup>68</sup> since it closes the loophole for Dirac particles near  $12\text{GeV}$ .) Note also that the generic constraints of Figure 6 also apply to other hypothetical particles since CDF and UA2 do not see any squarks, sleptons,  $W'$  or  $Z'$  up to masses significantly greater than  $M_{Z^0}$ . Thus, whatever the coupling boson is, it must be greater than  $M_{Z^0}$  which means the effective value for  $\sin^2 \phi_Z$  is  $< 1$ .

Furthermore, as Krauss<sup>67</sup> has emphasized, scalar particles such as sneutrinos interact like Dirac neutrinos so that the Kamiokande II and ionization experimental limits<sup>66</sup> also apply. Since asymmetric candidates are all Dirac particles, the restricted part of Figure 6b constrains asymmetric candidates where  $\Omega = 1$  is no longer required to follow the locus shown. Thus, it seems that whether the particle is matter-antimatter symmetric or not, it is required to have an interaction weaker than weak and/or have a mass greater than  $\sim 20\text{GeV}$ . Future dark matter searches should thus focus on more massive and more weakly interacting particles.

Also, as Dimopoulos<sup>63</sup> has emphasized, the next appealing crossing of  $\Omega = 1$  (see Figure 5) is  $\gtrsim 1\text{TeV}$  (but, in any case,  $\lesssim 340\text{TeV}$  from the unitarity bound<sup>65</sup>), which can be probed by SSC and LHC as well as by underground detectors. After the correct experimental constraints are taken into account, the favoured CDM particle candidate is now either a  $10^{-5}\text{eV}$  axion or a gaugino with a mass of many tens of  $\text{GeV}$ . Of course an HDM  $\nu_\tau$  with  $m_{\nu_\tau} \sim 20 \pm 10\text{eV}$  is still a fine candidate as long as galaxy formation proceeds by some mechanism other than adiabatic gaussian matter fluctuations.<sup>69,72</sup> This latter candidate becomes particularly attractive if recent hints from the gallium experiment<sup>73</sup> require the solution to the solar neutrino problem to have neutrino mixing with  $\nu_e - \nu_\mu$  mass scales of  $0.01$  to  $0.001\text{eV}$ , making multiple  $\text{eV}$  mass scales for  $\nu_\tau$  quite plausible from see-saw type models where  $m_{\nu_\tau} \sim m_{\nu_\mu} (\frac{M_{f_3}}{M_{f_2}})^2$  and  $M_{f_i}$  is an associated fermion mass for the  $i$ th generation. For example, if one uses the heavy quark masses,  $(\frac{M_t}{M_c})^2 \sim 10^4$ , so that  $\nu_\tau$  becomes ideal HDM. Such possibilities also may help late-time phase transition models for producing structure.<sup>72</sup>

The  $17\text{keV}$  reports are discussed in the appendix and are not dark matter candidates due to the instability requirement.

## Structure Formation

Perhaps the most outstanding problem in physical cosmology today is that of the formation of structure. Let us review the basic framework of structure formation in the universe. In particular, let us note that structure formation requires that density fluctuations grow. In order for this to occur,  $\rho_{\text{matter}}$  must be greater than  $\rho_{\text{radiation}}$ . If we

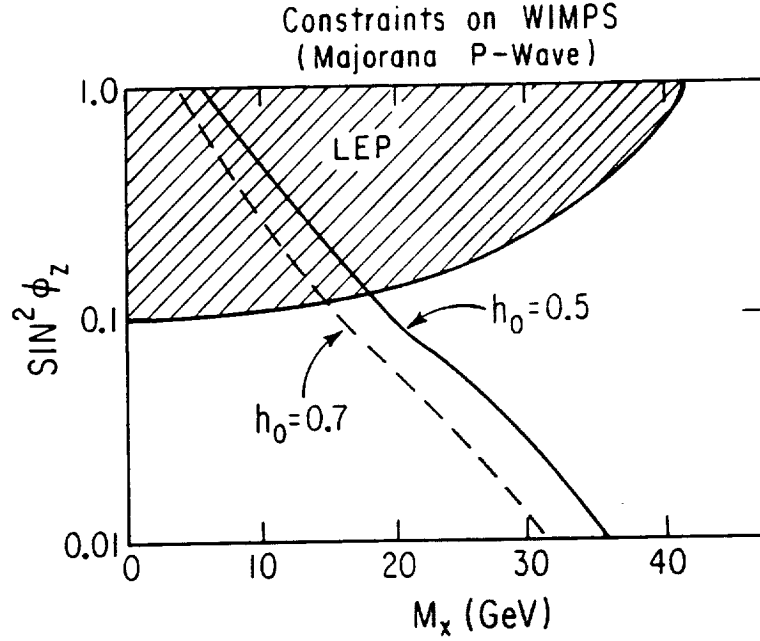


Figure 6a. Constraints on WIMPs of mass  $M_x$  versus  $\sin^2 \theta_z$ , the relative coupling to the  $Z^0$ . The constraints are shown assuming Majorana particle (p-wave interactions). The diagonal lines show the combinations of  $M_x$  and  $\sin^2 \theta_z$  that yield  $\Omega = 1$ . The cross-hatched region is what is ruled out by the current LEP results. Note that  $\Omega = 1$  with  $H_0 = 0.5$  is possible only if  $M_x \lesssim 20 \text{ GeV}$  and  $\sin^2 \theta_z < 0.1$ . This figure is revised from that of Ref. 68 using latest LEP results.

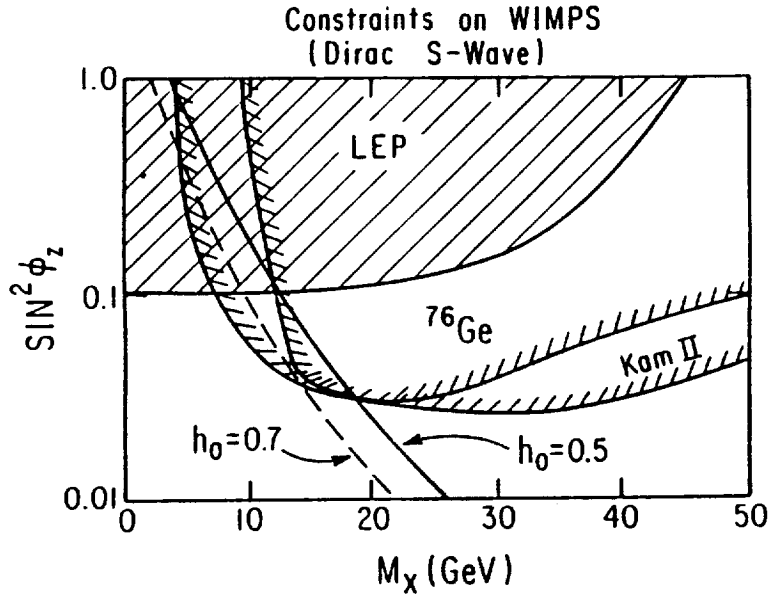


Figure 6b. This is the same as 6a but for Dirac particles (s-wave interactions). The  $^{76}\text{Ge}$  region is that ruled out by the Caldwell *et al.* double- $\beta$  decay style experiments. This figure is revised from that of Ref. 68 using latest LEP results and using new Kamiokande limits which closed a possible loophole near  $M_x \sim 10 \text{ GeV}$ . The current results require  $M_x \gtrsim 20 \text{ GeV}$  and  $\sin^2 \theta_x \lesssim 0.03$  for matter-antimatter symmetric particles and also exclude the entire cross-hatched region for asymmetric particle candidates.

define  $T_{eq}$  as the temperature where  $\rho_m = \rho_r$ , then for an  $\Omega = 1$  universe with  $h_0$  equal to 0.5, equality is approximately  $10^4$  times the present temperature  $T_0$ . The horizon mass at  $T_{eq}$  is  $\sim 5 \times 10^{16} (\frac{0.5}{h_0})^4 M$  which gives a present comoving scale of  $\sim 60 (\frac{0.5}{h_0})^2 Mpc$ . The recombination epoch  $T_{rec}$  for an  $\Omega = 1$  universe occurs slightly after matter domination. At  $T_{rec} \sim 1100 T_0$  baryon fluctuations begin to grow after recombination and the horizon mass at recombination is about  $10^{18} (\frac{0.5}{h_0}) M_\odot$  with a comoving scale of  $200 (\frac{0.5}{h_0}) Mpc$ . We also know that the fluctuations in the microwave background temperature at the time of recombination are less than a few parts in  $10^5$ .<sup>74</sup> Thus, in traditional models with primordial fluctuations existing prior to matter domination, growth begins at matter domination with the limits from  $\frac{\delta T}{T}$  forcing  $\frac{\delta \rho}{\rho}$  to be less than the order of  $10^{-4}$  since

$$\frac{\delta \rho_m}{\rho} \lesssim 3 \frac{\delta T}{T} \lesssim 10^{-4}.$$

Since small fluctuation  $\delta \rho$  grows linearly with  $1 + z$ , this would mean that fluctuations could reach the order of unity only at the present epoch. Non-linear growth, and thus true structure formation, does not begin until  $\frac{\delta \rho}{\rho}$  has reached unity (see Figure 7). Thus, in the standard model, the existence of objects at  $z > 1$  (see for example Gunn, Schneider, and Schmidt<sup>75</sup>) requires that there be fluctuations far larger than the average in order that these objects currently exist. As Efstathiou and Rees<sup>76</sup> point out, the gaussian fluctuation model for primordial fluctuations would not allow a large number of quasar-like objects to form at  $z \gtrsim 5$ .

All models for structure formation require at least two basic ingredients for that structure:

- (1) the matter,
- (2) the seeds.

In traditional models, the seeds are random fluctuations in the density field generated at the end of the GUT phase transition, presumably accompanying inflation.<sup>77,78</sup>

As mentioned in the previous section, the matter in any model of galaxy formation with  $\Omega = 1$  consists of normal baryonic matter with  $\Omega$  the order of 0.06 and some non-baryonic matter, either hot or cold, with  $\Omega$  the order of 0.94.<sup>79</sup>

The seeds which clump the matter to form objects may be divided into two broad categories (see Table 2) which can further be subdivided. The two broad categories would be (1) random gaussian seeds, presumably induced by quantum fluctuations at the end of a phase transition, and (2) topological defects produced in a vacuum phase transition. For the random gaussian seeds, the traditional assumption has been that the phase transition is the one associated with inflation.<sup>77,78</sup> However, it has been shown that similar kinds of fluctuations can also be generated in late-time phase transitions.<sup>15,16</sup> Similarly, for the topological defects, they could be formed either at the end of a GUT phase transition ( $\sim 10^{15} GeV$ ) or in some late-time transition.<sup>73,82,83</sup> In some sense this current division of random versus topological replaces the old division of adiabatic versus isothermal (or isocurvature). In fact, the current "random gaussian" are indeed "adiabatic" and the topological are isothermal and isocurvature. However, the latter have the new added feature of also being non-gaussian.

Let's note that all models for galaxy formation require new fundamental physics beyond the current particle standard model

$$SU_3 \times SU_2 \times U_1.$$

# Structure Growth

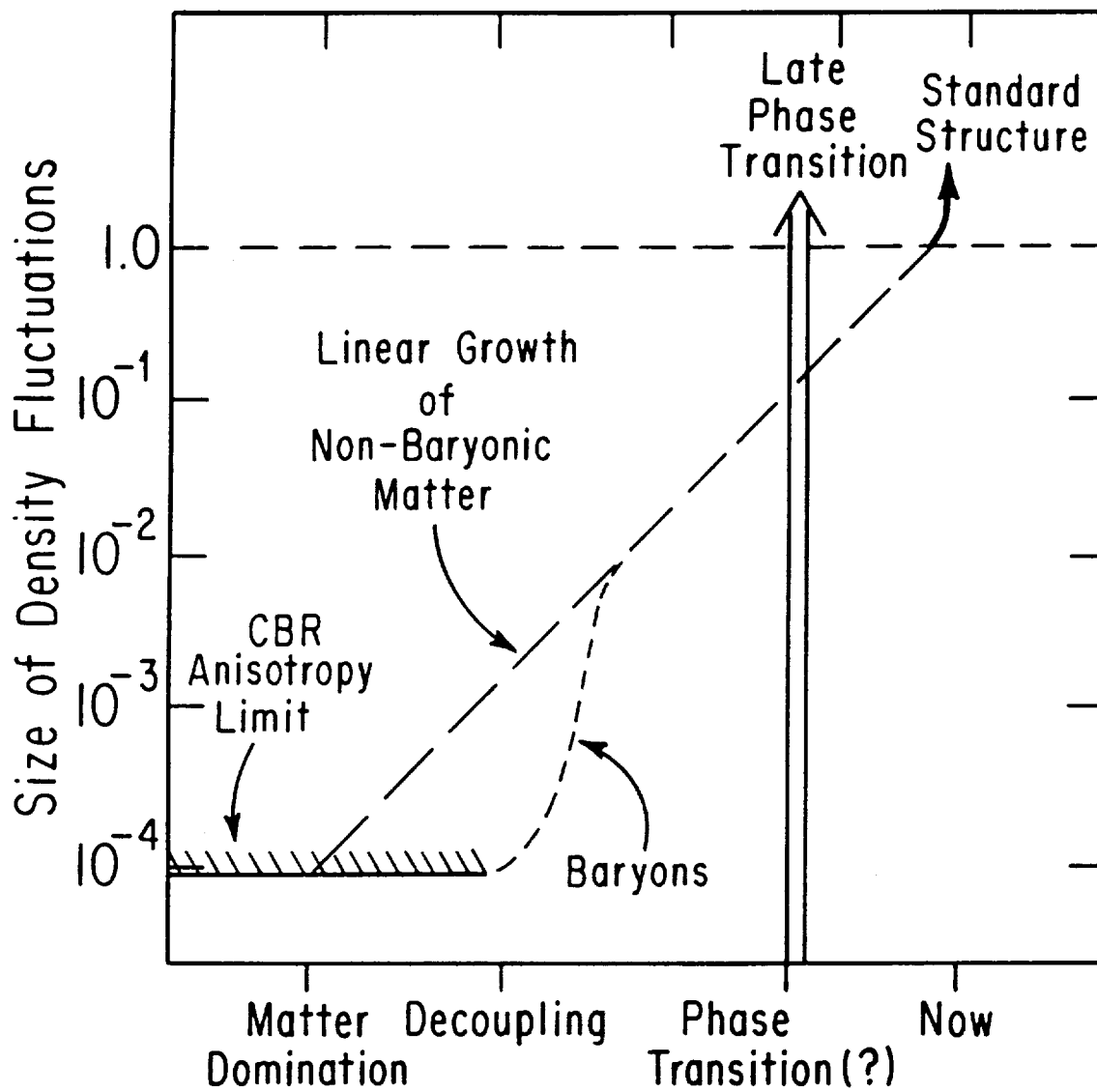


Figure 7. The growth of density fluctuations with the expansion of the universe. Note that LTPT can yield non-linear growth and thus structure at epochs much earlier than standard primordial models.

Table 2

SEEDS

I	RANDOM GAUSSIAN, Quantum Fluctuations
	A. End of Inflation
	B. LTPT
II	TOPOLOGICAL DEFECTS
	A. GUT
	B. LTPT

In particular, all non-baryonic dark matter, whether hot or cold, requires new physics, and similarly, all seeds, whether GUT scale or late-time and whether random gaussian or topological, require vacuum phase transitions. No model exists that does not invoke new physics. In fact, the existence of structure in the universe is one of the most important clues to the existence of physics beyond the standard model.

We should also note that not all combinations of seeds and matter are possible. For example, if one uses random gaussian seeds, then the non-baryonic matter must be cold, whereas if one uses topological seeds, the non-baryonic matter can be either hot or cold. One should also note that baryonic halos would require hot dark matter and hence topological seeds. Thus, searches for the dark baryons will also help constrain the non-baryonic candidates.

All current seed models require some form of vacuum phase transition. Thus, let us explore what possible phase transitions might occur (see Table 3). It should be noted in looking at Table 3 that of the three general classifications of cosmological phase transitions—the early, intermediate and late—the only ones that we absolutely know must have occurred are in the intermediate category when there is a horizon problem, namely that the horizon at the time of that transition is too small to generate galactic sized structure, and yet, the transition is not accompanied by significant inflation. The traditional early transitions have been used in the past because, while their horizon is small, inflation can amplify the effects to large scales. The other option is that of a late-time transition, where the universe waits until the horizon is sufficiently large that the physics of the phase transition directly yields the structures without having to use inflation to avoid the horizon problem.

## Potential Observations to be Explained

In the last couple of years there have been a number of observations affecting galaxy formation and large-scale structure that have been a potential problem for traditional models which invoked early random gaussian fluctuations. However, because each of these observations is new and has not stood the test of time, in this discussion we refer to these as potential observations. In particular, many of the advocates of gaussian fluctuations and cold dark matter have tried to argue that these observations are statistical flukes that have yet to be established. Obviously, if these potential observations continue to hold up and are verified and are shown to be ubiquitous rather than statistical rareties, then the traditional models are in serious trouble. Table 4 summarizes these potential observations. Perhaps the most potentially damning would be observations of microwave anisotropies  $\frac{\Delta T}{T}$  at levels significantly below  $10^{-5}$ . However, at the present time, observations of small scale anisotropy are at the level of a couple times  $10^{-5}$ . Observations on angular scales of degrees or more are also approaching a few  $10^{-5}$ . As this paper is being written, the measurements have not yet reached the point of ruling out the model of random fluctuations. However, as noted by Smoot,<sup>74</sup> within the not too distant future, COBE may be able to achieve limits as low as  $3 \times 10^{-6}$  on scales of a few degrees and larger, and antarctic studies may also push to similar levels on somewhat smaller scales, as might the balloon studies of Meyer *et al.* at MIT.

The next observation that can be a potential problem for traditional models is the existence of structures with scales greater than the order of  $100Mpc$ . In particular, the great wall observed by Geller and Huchra<sup>84</sup> shows that there is at least one such wall in the universe. The observations of Broadhurst *et al.*<sup>85,86</sup> show evidence for a multiplicity of

Table 3

VACUUM PHASE TRANSITIONS

EARLY	(Small horizon but inflation)
	$\sim 10^{19} GeV$ - T.O.E.
	$\sim 10^{16} GeV$ - GUT
INTERMEDIATE	(Known to occur but horizon problem)
	$\sim 10^2 GeV$ - Electroweak
	$\sim 1 GeV$ - QCD
LATE	(Horizon large)
	$\sim 10^{-2} eV$ - Family symmetries, etc.



Table 4

POTENTIAL OBSERVATIONS

1.  $\frac{\delta T}{T} \lesssim 10^{-5}$
2. Structures  $\gtrsim 100 Mpc$
3. Large coherent velocity flows
4. Objects existing at  $z \gtrsim 5$
5. Large cluster - cluster correlations

such great walls with the characteristic spacing comparable to the size of the Geller-Huchra wall itself. While much debate has been made about whether or not the multiple walls of Broadhurst *et al.* are periodic or quasi-periodic, it does seem clear from their observations, as well as the work reported by Szalay,<sup>87</sup> that there is significant structure in the universe on scales of  $\sim 100Mpc$ . This is thoroughly supported by the large coherent velocity flows where the Seven Samurai<sup>88</sup> and others have found evidence for the existence of an object they call the “Great Attractor” towards which the Virgo cluster and the Hydro-Centaurus cluster all seem to be flowing with a velocity  $\sim 600km/sec$ . This again seems to indicate evidence of structures on the scales of at least  $60Mpc$ .

Perhaps most constraining of the traditional astronomical measurements is the existence of objects at very large redshifts. In particular, Schneider, Schmidt and Gunn<sup>75</sup> have found a quasar with a redshift of 4.73 (and they have privately reported one at 4.9). As Efstathiou and Rees<sup>76</sup> have noted, if such objects are ubiquitous, this would be fatal for primordial gaussian fluctuation models. Similarly, if one ever finds a quasar-type object at much larger redshifts, that would also be fatal.

Another potentially fatal observation for gaussian fluctuation models comes from the work of Bahcall and Soneira,<sup>89</sup> and Klypin and Khlopov<sup>90</sup> where they find that clusters of galaxies seem to be more strongly correlated with each other than galaxies are correlated with each other. While Primack and Dekel<sup>91</sup> have warned of the dangers of projection effects on such observations, it seems difficult to understand how projection effects would give the fractal-like behavior.<sup>92</sup> Furthermore, the southern hemisphere work of Huchra<sup>93</sup> also seems to support high cluster correlations. Most recently Vandenburg and West<sup>94</sup> have also found similar correlations for the CD galaxies observed at cluster centers. These CD’s should not have the projection effect problems because redshifts are known. Even Primack and Dekel now acknowledge that there seems to be some excess in cluster correlations. If such large correlations turn out to be real, they too cannot be easily explained in the gaussian model, and, as Szalay and Schramm<sup>93</sup> note, they seem to be best fit by some sort of fractal-like pattern, as one might get from topological defects induced by a phase transition.

## Late-Time Transitions

By late-time transition we will mean any non-linear growth occurring shortly after recombination. As mentioned above, such non-linear growth can be related either to a gaussian pattern or to a topological pattern such as walls, strings or textures. It is also possible that some normal random gaussian pattern from the very early universe could be triggered to undergo non-linear growth by some sort of phase transition or related phenomenon occurring after recombination. An example of this latter case would be the neutrino flypaper model of Fuller and Schramm.<sup>95</sup>

In general we will see that these late-time transitions can give the smallest possible  $\frac{\Delta T}{T}$  for a given size structure. They can produce non-gaussian structural patterns, fractal-like with large velocity flows. It might be noted that the co-moving horizon at the time of the transition is not too different than the scale associated with the largest structures observed. No model of primordial fluctuations naturally imbeds this horizon scale onto the structural pattern. If some non-linear growth is associated with the patterns, the horizon scale can be imposed on the structure.

Another very dramatic advantage of late-time transitions, illustrated in Figure 7, is

that it can produce structure with  $\frac{\delta\rho}{\rho} \gtrsim 1$  at  $z \gtrsim 10$ . Thus, one could have significant structure and a significant number of objects at high redshift, which is a problem in any normal model with the seeds forming prior to recombination.

Let us now explore the possible physics that might give rise to a late-time transition, that is, a transition with a critical temperature between 0.001eV and 1eV. It might be noted that in some sense it is a “hierarchy” rather than a “fine-tuning” problem to obtain a transition in this temperature range. We are trying to find a small mass scale somewhat analogous to how one would like to find the mass scale of the electron, or, for that matter, the  $Z^0$  boson, when the natural mass scales to the problem are closer to  $10^{19}GeV$ , as in superstring models, or to 0. The hierarchy problem of trying to find the intermediate scale of the electroweak interaction of somewhere between the quark-lepton scale and the GUT or Planck scale has traditionally been approached with either a supersymmetric solution or a dynamical solution (“technicolor”). This supersymmetric solution, in some sense, is analogous to the model proposed in the appendix of Hill, Schramm and Fry,<sup>72</sup> denoted as HSF, which is an adaptation of the Hill-Ross<sup>96</sup> mechanism. A dynamical solution which has been proposed by Dimopoulos<sup>97</sup> involves a shadow  $SU3$ . The scale of a physics that might be associated with an HSF mechanism was relating to the MSW mixing solution to the solar neutrino problem.

The MSW<sup>98,99</sup> mixing solution to the solar neutrino problem is achieved if the neutrino mass difference squared,  $\delta m^2$ , is of the order of  $10^{-4}$  to  $10^{-7}eV^2$ , or, in other words, neutrino masses of the order of a fraction of an electron volt. If we assume, following HSF, that the neutrino masses are generated by a pseudo-Nambu-Goldstone boson mechanism with mass

$$m_\phi \sim \frac{m_\nu^2}{f}$$

and with a transition occurring at  $T_{crit} \sim m_\nu$ , and if we further assume that the coupling  $f$  is related to the GUT scale, since we want to imbed this in some sort of unified theory, then the Compton wavelength  $\lambda_\phi \sim 1Mpc$ , in other words, a galactic scale. The density of the  $\phi$  field at the time of the transition is the order of the cosmological density, in other words,

$$\frac{\rho_\phi}{\rho} \sim 1.$$

(Note that this is natural for phase transitions, whereas the requirement for primordial transitions to have small fluctuations, as inflation requires, is a fine tuning requirement.) Furthermore, the average spacing of the nucleation sights,  $L$ , can be estimated from Coleman’s theory on spontaneous nucleation to yield spacings today that are interesting:

$$\frac{R_H}{L} \sim \text{Log}\left(\frac{M_p}{T_{crit}}\right)$$

$$L_{co} \equiv L(1 + z_{crit}) = \frac{6000}{\text{Log}\left(\frac{M_p}{T_{crit}}\right)} \left(\frac{0.5}{h_0}\right) (1 + z_c)^{-1/2}$$

where  $z_{crit} = z_c = (T_{crit}/T_0 - 1)$ ,

$R_H$  is the horizon radius at  $z'_c$  and  $M_p \sim 10^{19} \text{ GeV}$ . This yields for  $T_{crit} \sim 10^{-2} \text{ eV to } 10^{-3} \text{ eV} / L_{co(moving)} \sim 40 \text{ to } 140 \text{ Mpc}$ .

As we mentioned previously, recent impetus for new physics at this energy scale has come from the SAGE experiment which detects neutrinos from the PP chain in the sun. The previous solar neutrino experiments, the chlorine and the Kamiokande experiments, are mainly sensitive to the rare  $^8\text{B}$  branch of the solar energy generating reactions. It is well established that the  $^8\text{B}$  experiments have seen fluxes at levels somewhat below theoretical predictions.<sup>100</sup> However, there has always been the worry that the  $^8\text{B}$  channel may be suppressed due to astrophysical effects since its yield is very temperature sensitive. However, the PP chain that produces the neutrinos to be detected by SAGE must work if the sun is burning by fusion. Thus, the report<sup>101</sup> of no significant counts above background after five months of running the gallium experiment when they expected nineteen counts for the standard model implies that something is happening to the neutrinos on their way between emission and arrival at earth. (Or, that something is wrong with the detector, such as  $^{71}\text{Ge}$  produced by  $\nu$ -capture because it starts as an ion may have different chemistry than neutral  $^{71}\text{Ge}$ .) Of course, the present results are very preliminary. Questions with regard to estimates of background, counting efficiencies, systematics, statistics, etc., remain, but the tantalizing hint that the  $\nu_e$ 's mixed into some other species of neutrino on their way out of the sun is certainly exciting. The final state of this experiment will not be known for several years. The similar gallium experiment operated by the GALLEX collaboration in the Grand Sasso Tunnel in Italy is also beginning to run but so far has had some background problems. The GALLEX chemistry may be somewhat cleaner and we will thus have an independent check on SAGE. Furthermore, both of these gallium experiments will be calibrated using  $^{51}\text{Cr}$  sources of  $\text{MeV}$  neutrinos. Thus, one will have a true check of their counting efficiencies, etc., and both of these experiments will run for a long-enough time that the statistics will reach significant levels. If the neutrinos really are mixing on their way out of the sun, then the MSW solution is probably valid and we are in the realm discussed above.

It might also be noted that a simple application of the Gell-Mann-Ramond-Slansky see-saw model<sup>102</sup> for neutrino masses yields some interesting implications. If we assume that there is a mass hierarchy in the neutrinos with the electron neutrino having negligible mass, the  $\mu$  the intermediate mass and the  $\tau$  the heaviest, and we assume that the mixing of the  $\nu_e$  in the sun goes to its nearest neighbor family, the  $\nu_\mu$ , then the  $\nu_\mu$  is carrying most of the mass of the MSW  $\delta_m^2$ . The see-saw mechanism argues that

$$m_{\nu_i} \sim \frac{m_{f_i}^2}{M}$$

for a given family, or, in other words,

$$m_{\nu_\tau} \sim m_{\nu_\mu} \left( \frac{m_{f_3}}{m_{f_2}} \right)^2.$$

If we use lepton masses for the fermion masses, this yields a  $\nu_\tau$  mass in the neighborhood of a few  $\text{eV}$ . However, if we use heavy quark masses, then, since the top quark mass is  $\gtrsim 100$  times that of the charm quark, this yields  $\nu_\tau$  masses in the neighborhood of 10 to 100  $\text{eV}$ , making it perfect hot dark matter. It might also be noted that the see-saw mass scale,  $M$ , in this picture, ends up being the order of  $10^9$  to  $10^{12} \text{ GeV}$ , which happens to

be the only window allowed for the DFS-axion<sup>103</sup> scale. It might further be noted that if the non-baryonic dark matter is indeed the  $\tau$  neutrino, then one is required to dismiss primordial gaussian fluctuations.

Note that even if the MSW mixing is  $\nu_e - \nu_\tau$ , the LTPT possibility is still there, but then all neutrinos would be light and could not serve as HDM. It is interesting that in this latter case the see-saw  $M$  is the GUT scale.

## Structure from LTPT

LTPT can produce vacuum fluctuations of the random gaussian character just as could be generated at the end of inflation.<sup>76</sup> However, as emphasized in references 80 and 81, these structures will have a quantum scale that is the order of a galaxy size, and the bosons associated with the fluctuations might even serve as the dark matter of the universe.

The other alternative for LTPT is to produce topological structures. Just as early universe phase transitions can produce strings and/or textures, LTPT can also produce such objects. Furthermore, LTPT can produce walls which are a problem for primordial phase transitions. However, there is a problem for some walls, depending on the nature of the interaction potential. LTPT that have a  $\lambda\phi^4$  potential will end up with one wall dominating as was demonstrated in references 104, 105 and 106. However, this problem of one wall dominating can be surmounted in a variety of ways which have varying degrees of attractiveness, depending on the eyes of the beholder. For example, in the HSF phase transition, the walls are sine-Gordon rather than  $\lambda\phi^4$ . As Widrow has shown,<sup>107</sup> the sine-Gordon walls can yield "bags" of wall or "balls" of wall which survive several expansion times. These bags or balls can then serve as seeds in galaxy formation, and thus, it is their amplitude that becomes a deciding factor for  $\frac{\Delta T}{T}$  limits as opposed to the energy scale of the infinite walls which can be made quite small. This latter point was emphasized by Hill, Schramm and Widrow.<sup>82</sup> Another way of avoiding single wall dominance is the decaying wall model of Kawano<sup>108</sup> where the walls serve as seeds and then decay away. It is also possible to escape one-wall domination with a large number of minima in the potential. Perhaps the most dramatic way of escaping one-wall domination, thus keeping a network of walls, as shown in Figure 8A, is if the walls have friction with the ambient medium, whether it be neutrinos or the remaining baryonic and/or non-baryonic matter in the universe.<sup>109</sup> Alessandro Massarotti has shown that friction can in many reasonable cases slow the walls down sufficiently that they do not evolve to the one-ball domination situation. In this case, one retains a complex network with  $L$  for the wall being much less than the horizon size.

It might be noted that long walls gravitationally repel rather than attract,<sup>110,111</sup> whereas balls of wall are attractive seeds. Thus, a combined network of balls and slowed-down long walls can yield a complex structure which may be even of a fractal character in agreement with the claims of Schramm and Szalay<sup>92</sup> from cluster correlations.

In addition to walls, LTPT can also produce textures<sup>112</sup> or non-topological solitons.<sup>113</sup> In these latter cases, or with the bags of wall dominating, one will have networks more closely resembling Figure 8B and Figure 8A. It should be noted that the parameters  $L$  and  $\delta$  and the nature of the structures generated are dependent on the model for the LTPT. It should also be noted that questions of the detailed physics of imbedding the LTPT into some larger GUT or TOE are dependent on the unification model. HSF have shown that a reasonable toy model can be constructed which can give a phase transition.

# WALL NETWORK

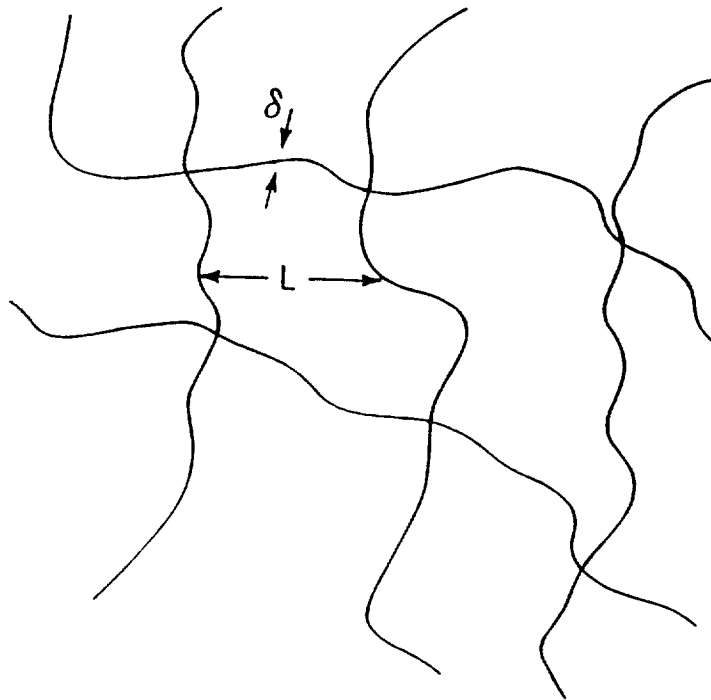


Figure 8a. A generic wall network defining the wall thickness  $\delta$  and the characteristic spacing of structure  $L$ .

## SEED NETWORK (Bag, balls-of-wall, Textures, etc)

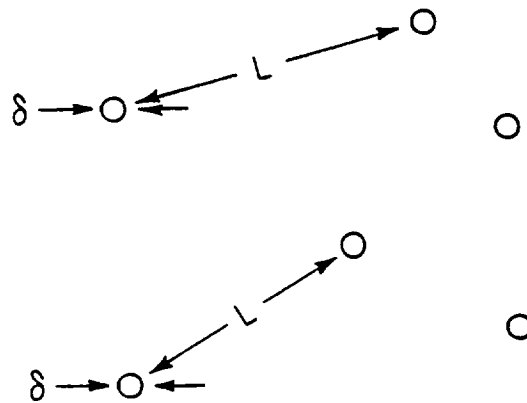


Figure 8b. A generic network for seed generation with seed size  $\delta$  and seed separation  $L$ .

These phase transitions in many ways are quite analogous to the axion-producing phase transition which has a coupling at a scale near to the order of  $10^{11} GeV$ , far above the QCD phase transition scale of the order of  $GeV$ . And like the axion, the particle involved in the LTPT of HSF has a pseudo-Nambu-Goldstone boson. However, instead of being related to the strong interaction and quarks, in the LTPT case it is related to the neutrinos and probably to family symmetry.

Generating seeds at an LTPT might be advantageous for producing the multiple walls of Broadhurst *et al.*<sup>85,86</sup> In particular, Icke and Weygaert,<sup>114</sup> and Coles<sup>115</sup> have independently demonstrated that the phenomenological Voronoi tessalations of the intersection of expanding rarefaction shells give a very good fit to large scale structure if the nodes of these tessalations are fit to the Abell clusters. In particular, they note that one gets quasi-periodic walls at  $\sim 130 Mpc$  with cluster correlation functions that are quite strong and follow the fractal behavior of Schramm and Szalay. However, the seed distribution required to give this tessalation causes a conflict with the microwave background radiation, if the seeds are generated prior to the decoupling. However, an LTPT could remedy that. Similarly, an LTPT can provide the seeds to enable hot dark matter to work as a galaxy formation model (see, for example, reference 116.) It might be noted that the typical bag of wall can easily yield a galaxy or a quasar-forming seed.

We can estimate the mass associated with a wall in the following way:

Let  $\sigma$  = energy density per unit area, that is:

$$\sigma \equiv \rho_w \delta \sim \frac{4 \times 10^{-5} s \rho_0 \delta}{h_0^2} (1 + z_c)^4$$

where

$\delta \equiv$  thickness

$\rho_w$  = density of the wall

$$s = \frac{\rho_w}{\rho_r} \text{ at } z_c$$

$$\rho_0 = 3 \times 10^{11} h_0^2 M_\odot / Mpc^3$$

then

$$M_w \sim \pi \sigma L^2 \sim 3 \times 10^7 (1 + z_c)^4 \left( \frac{\delta}{Mpc} \right) \left( \frac{L}{Mpc} \right)^2 M_\odot$$

and for stable walls

$$\Omega_w(z) \sim \frac{3}{4} \frac{\sigma}{\rho_0 L (1 + z)^3}.$$

Note that  $\Omega_w$  at the present epoch, can be the order of unity. Wall domination can occur at present epoch if

$$z_c \gtrsim 11 \left( \frac{L}{\delta} \right)^{\frac{1}{4}} \left( \frac{h \sigma^2}{s} \right)^{\frac{1}{4}} - 1$$

for stable walls. It might be noted that if wall domination occurs at the present epoch, as long as there are multiple walls, rather than just one wall dominating, one has the interesting situation where the expansion of the universe is no longer following the normal matter-dominated relationship, and, in particular, one can achieve ages greater than  $\frac{1}{h_0}$ . Such a situation may be a solution to the age-Hubble constant problem if  $h_0$  is ever shown to be greater than 0.7.

It might also be noted for topological structure generated by LTPT that the structure is relatively independent of whether the non-baryonic dark matter is hot or cold.

## Microwave Anisotropies

Since LTPTs provide no fluctuations on the surface of last scattering, all fluctuations from the microwave background must be due to the differential redshift-blueshift non-cancellation due to a changing potential in the transparent medium or due to scattering of the microwave photons off of moving objects. One can estimate the potential change due to the  $\phi$  field itself generated in the phase transition and by the dynamic motion of the structures and the Doppler shift thereby produced. One can also do the classical Rees-Sciama and Sachs-Wolf calculations for the  $\frac{\Delta T}{T}$  generated by existing objects.<sup>118,119</sup> We can estimate its effects roughly in the following way: The static effects will dimensionally go as

$$\frac{\delta T}{T} \sim \Omega_w \left( \frac{L}{R_H} \right)^2 \sim G\sigma L.$$

The time-changing effects can be estimated by multiplying the static effect by  $\frac{V}{c}$ . While different people remember different formulations of these things, one can show that because of the nature of walls and other topological systems, the effects can be reduced to the form  $G\sigma L$  times  $\frac{V}{c}$  or  $\frac{V^2}{c^2}$ . Since any walls or topological seeds we ever see must be moving with  $V < c$ , the dominant effect will in general go like  $G\sigma L$ , which can be shown to yield the result:

$$\frac{\delta T}{T} \sim 10^{-8} \left( \frac{1+z_c}{10} \right)^4 s \left( \frac{\delta}{Mpc} \right) \left( \frac{L}{Mpc} \right)$$

$$\sim 10^{-6} \text{ for } L \sim 100 Mpc, \delta \sim 1 Mpc, z_c \sim 10.$$

Note that this yields  $\frac{\delta T}{T} \sim 10^{-6}$  even for an  $L$  of  $100 Mpc$ . The distribution, however, of these fluctuations depends very much on the detailed topological nature of the structures produced. In particular, Turner, Watkins and Widrow<sup>119</sup> have shown that balls of wall tend to produce spikes very similar in nature to the spikes that textures produce. A general formalism showing the wide range of structures in non-gaussian microwave background fluctuations has been developed by Goetz and Noetzold.<sup>120</sup>

In general, one can see that if structure of size  $L$  is generated by a late-time phase transition, and  $L$  is the maximum size of structure produced in that transition, then the late-time transition does give the minimum  $\frac{\delta T}{T}$  for that structure. Of course the question is what is the characteristic size  $L$  of structure generated in a transition. For  $\lambda\phi^4$  structures,  $L$  goes to the horizon size, in which case  $\frac{\delta T}{T}$  gets larger than current observational limits. However, as mentioned above, many other possibilities can be generated in LTPT, with



L at present being a somewhat freely adjustable parameter, depending on the model, the amount of friction, the decay of the walls, etc.

## Conclusion

In these lectures we have seen that the basic big bang model is in excellent shape with the recent collider results helping to confirm it in the same way they've helped with  $SU_3 \times SU_2 \times U_1$ . We've also seen that the prediction that the bulk of the matter in the universe is in the form of some exotic non-baryonic species obviously remains to be confirmed. Furthermore, we've examined the current problems of generating structure in the universe, mentioning the traditional primordial scenarios as well as the new exotic idea of a late-time phase transition.

All in all, we've seen that cosmology is tremendously active with new data coming from both astronomical and particle physics techniques. We've also seen that some of the best current indications for particle physics beyond  $SU_3 \times SU_2 \times U_1$  are coming from astrophysical arguments.

## Acknowledgments

I would like to thank my recent collaborators, Xiaochun Luo, Chris Hill, John Ellis, Savas Dimopoulos, Gary Steigman, Keith Olive, Michael Turner, Rocky Kolb and Terry Walker, for many useful discussions. I would also like to acknowledge useful communications with Denys Wilkinson, Walter Mampe, Jim Byrne and Stuart Friedman on the neutron lifetime and with Bernard Pagel and Jay Gallagher on helium abundances and Alessandro Massarotti and Neil Turok on phase transitions. This work was supported in part by NSF grant AST 88-22595, by NASA grant NAGW 1321 and by DoE grant DE-FGO2-91ER40606 at the University of Chicago and by the DoE and NASA grant NAGW 1340 at the NASA/Fermilab Astrophysics Center.

## Appendix I - The 17 keV "Thing"

A few years ago Simpson<sup>122</sup> reported the possible existence of a 17 keV mass state mixing at the 1% level with an electron neutrino during tritium decay. Although initially met with some skepticism, it could not be trivially dismissed. Recently, attention has been refocused on this object as a result of a series of other nuclear  $\beta$  decay experiments. In particular, Norman<sup>123</sup> and his collaborators reported the existence of a similar 17 keV component at the 1% mixing level in carbon-14 decay, and, most dramatically, Hime and Jelly<sup>124</sup> reported a similar 17 keV mass mixing at the 1% level in sulphur-35 decay. While this is all still somewhat preliminary, and there are worries that there could be some sort of non-obvious instrumental effect occurring (for example, it might be noted that the effect has been seen only using solid state detectors rather than magnetic spectrometers<sup>125</sup>), nonetheless, the possibility that an electron neutrino mixes at 1% level with something that has a 17 keV mass has raised much excitement.

There's a problem, however, in what this object can be. Numerous papers have been written discussing the problems and trying to come up with exotic models that might be able to fit it.<sup>126</sup> Basic problems are as follows:

1. It is well known that there are only three families of neutrinos from the LEP experiments. Thus, if the 17 keV object is another neutrino, it must be either the  $\mu$  or the

$\tau$ . But limits on  $\nu_e - \nu_\mu$  mixing show that it is much less than 1%, so the only neutrino possibility is  $\nu_\tau$ .

2. It is also known that neutrinoless double  $\beta$  decay has not been seen. This limits the majorana mass for the electron neutrino to be less than approximately 1 electron volt. Since the 17 keV mass times the 1% mixing yields a mass of 170eV, this particle cannot have a majorana mass in the normal interpretation. (The exotic option of having the majorana mass term be a combination of  $\nu_e$  ( $C\bar{\nu}_\mu + S\bar{\nu}_\tau$ ) and thus maximally violate the CP cannot be easily excluded.<sup>127</sup>)

3. If the 17 keV does not have a majorana mass, then it must have a Dirac mass if it is a neutrino.

4. If it has a Dirac mass, then big bang nucleosynthesis would count the spin-flip component as an extra neutrino state and, as we saw earlier, big bang nucleosynthesis does not allow more than a total of 3.4 neutrinos, thus excluding any of the normal neutrinos from having right-handed components that interact with normal neutrino-like interactions.

5. The remaining option would be that this right-handed component must interact much more weakly than normal left-handed neutrinos. Olive *et al.*<sup>128</sup> showed that the nucleosynthesis limits on exotic neutrinos can be surmounted if those exotic neutrinos interact much more weakly than normal neutrinos. Thus, the right-handed component would have to couple to a  $\mu c Z'$  with a mass greater than about  $\sim TeV$ .

6. This leads to a dilemma in another astrophysical area. If the right-handed neutrino is so weakly coupled, then it would have freely escaped from SN 1987A. Gandhi and Burrows<sup>129</sup> have argued that right-handed neutrinos with masses greater than 14 keV are excluded (a recent numerical error in their calculation may push their limit up to 28 keV). While this limit at first appears only marginal, it should be noted that calculations treating neutrino processes in more detail (including neutrino bremsstrahlung induced spin-flip which Turner<sup>130</sup> has shown will significantly enhance production) as well as the slightly higher temperatures encountered in the course of other supernova collapse calculations (see, for example, Mayle, Wilson, Schramm<sup>131</sup>) appear to strengthen this limit and seem to push it down significantly below  $\sim 10 keV$ , thus severely constraining the existence of a right-handed Dirac neutrino. Furthermore, one can't have it both ways. If the neutrino is sufficiently weak to escape the cosmological bound, then it makes it easier for it to get out of the supernova and make the time-scale of the neutrino burst in SN 1987A shorter than what was observed to be. The basic physics in the time-scale is simply that the ten second duration of the neutrino burst (see review by Truran and Schramm<sup>132</sup>) requires that neutrinos diffuse out rather than freely stream out. If any component of the neutrinos is able to stream out freely, then there would be a leakage of energy out of the core and the duration of the neutrino pulse would be much less. Similar arguments to this were used to set limits on axions and any other exotic particles that might have been produced in the collapsing core.

7. From the cosmological constraints on  $\Omega$ , any 17 keV neutrino would have to be unstable, since a 17 keV stable neutrino would yield an  $\Omega$  of approximately 200, which would have led to the Big Crunch many eons ago.

8. If the neutrino were unstable, it would have to be sufficiently unstable that it did not escape the supernova core, and it could not decay to photons since no gamma rays were seen to accompany the neutrinos from SN 1987A. Thus, in the  $10^5$  year lifetime of transit from 1987A to the solar system, the neutrinos did not produce significant radiative

products.

9. This further constrains any model since an invisible decay that did not produce photons would require some other new particles such as a majoron, which then should have been counted in big bang nucleosynthesis and, as we've already seen, that limit seems to be quite formidable.

All in all, this seems to mean that whatever the 17keV thing is, it is not a neutrino in any normal sense of the word.

## References

1. J. Mather *et al.*, COBE preprint, Goddard Space Flight Center (1990) and *Astrophys. J.*, in press (1991).
2. D.N. Schramm, in *Proc. of the 1990 Rencontre de Physique at La Thuile, March 1990*, ed. M. Greco (in press, 1991);  
D.N. Schramm, in *Proc. 1990 Nobel Symposium, Gräfvållen, Sweden* (in press, 1991);  
D.N. Schramm, in *Proc. of the Moriond Astrophysics Meeting on Particle Astrophysics at Les Arcs, Savoie, France, March 1990*, ed. T. Piran (in press, 1991);  
D.N. Schramm, in *Z<sup>0</sup> Physics - Cargèse 1990*, eds. M. Levy *et al.* (Plenum, London, 1991).
3. R.A. Alpher, H. Bethe and G. Gamow, *Phys. Rev.* **73** (1948) 803.
4. R.A. Alpher, J.W. Follin and R.C. Herman, *Phys. Rev.* **92** (1953) 1347.
5. R. Taylor and F. Hoyle, F., *Nature* **203** (1964) 1108.
6. P.J.E. Peebles, *Phys. Rev. Lett.* **16** (1966) 410.
7. P. Wagoner, W.A. Fowler and F. Hoyle, F., *Astrophys. J.* **148** (1967) 3.
8. D.N. Schramm and R.V. Wagoner, *Ann. Rev. of Nuc. Sci.* **27** (1977) 37;  
K. Olive, D.N. Schramm, G. Steigman, M. Turner and J. Yang, *Astrophys. J.* **246** (1981) 557;  
A. Boesgaard and G. Steigman, *Ann. Rev. of Astron. and Astrophys.* **23** (1985) 319.
9. J. Yang, M. Turner, G. Steigman, D.N. Schramm and K. Olive, *Astrophys. J.* **281** (1984) 493.
10. K. Olive, D.N. Schramm, G. Steigman and T. Walker, *Phys. Lett. B.* **236** (1990) 454.
11. T. Walker, G. Steigman, D.N. Schramm, H.-S. Kang and K. Olive, *Astrophys. J.* (submitted, 1991).
12. L. Kawano, D.N. Schramm and G. Steigman, *Astrophys. J.* **327** (1988) 750.
13. J.W. Truran, A.G. Cameron and A. Gilbert, *Canadian Journal of Physics* **44** (1966) 563.
14. L. Krauss and P. Romanelli, P., Yale University preprint (1990) and *Ap.J.*, in press (1991).
15. W. Fowler, J. Greenstein and F. Hoyle, *Geophys. J.R.A.S.* **6** (1962) 6.
16. C. Ryter, H. Reeves, E. Gradstajn and J. Audouze, *Astron. and Astrophys.* **8** (1970) 389.
17. J. Geiss, and H. Reeves, *Astron. and Astrophys.* **18** (1971) 126;  
B. Black, *Nature* **234** (1973) 148.

18. J. Rogerson and D. York, *Astrophys. J.* **186** (1973) L95.
19. H. Reeves, J. Audouze, W.A. Fowler and D.N. Schramm, *Astrophys. J.* **179** (1973) 909.
20. R. Epstein, J. Lattimer and D.N. Schramm, *Nature* **263** (1976) 198.
21. J.R. Gott, III, J. Gunn, D.N. Schramm B.M. Tinsley, *Astrophys. J.* **194** (1974) 543.
22. R.T. Rood, G. Steigman and B.M. Tinsley, *Astrophys. J.* **207** (1976) L57.
23. J. Yang, D.N. Schramm, G. Steigman and R.T. Rood, *Astrophys. J.* **227** (1979) 697.
24. T. Wilson, R.T. Rood and T. Bania, in *Proc. of the ESO Workshop on Primordial Healing*, eds. P. Shaver and D. Knuth (European Southern Observatory, Garching, 1983).
25. M.R. Hartoog, *Ap.J.* **231** (1979) 161;  
J. Ostriker and D.N. Schramm, Fermilab/Princeton preprint (1991).
26. J. Spite and F. Spite, *Astron. and Astrophys.* **115** (1982) 357;  
R. Rebolo, P. Molnaro and J. Beckman, *Astron. and Astrophys.* **192** (1988) 192.  
L. Hobbs and C. Pilachowski, *Astrophys. J.* **326** (1988) L23.
27. G. Steigman, D.N. Schramm and J. Gunn, *Phys. Lett.* **66B** (1977) 202.
28. V.F. Schwartzman, *JETP Letters* **9** (1969) 184;  
P.J.E. Peebles, *Physical Cosmology* (Princeton University Press, 1971).
29. D.N. Schramm and L. Kawano, *Nuc. Inst. and Methods A* **284** (1989) 84.
30. B. Pagel, *Proc. of 1989 Rencontres de Moriond* (1990).
31. R. Scherrer, J. Applegate and C. Hogan, *Phys. Rev. D* **35** (1987) 1151;  
C. Alcock, G. Fuller and G. Mathews, *Astrophys. J.* **320** (1987) 439;  
W.A. Fowler and R. Malaney, *Astrophys. J.* **333** (1988) 14.
32. H. Kurki-Suonio, R. Matzner, K. Olive and D.N. Schramm, *Astrophys. J.* **353** (1990) 406.
33. M. Peimbert and S. Torres-Peimbert, *Astrophys. J.* **193** (1974) 327.
34. B. Pagel, in *ESO/CERN Proc.* (1990).
35. G. Steigman, D.N. Schramm and J. Gallagher, *Comments on Astrophys.* **14** (1989) 97.
36. W. Mampe, P. Ageron, C. Bates, J.M. Pendlebury and A. Steyerl, *Phys. Rev. Lett.* **63** (1989) 593.
37. J. Byrne, *et al.*, *Phys. Lett. B*, submitted (1991).
38. H. Abele, M. Arnould, H.A. Borel, J. Dohner, D. Dubbers, S. Freedman, J. Last and I. Reichert, in *Proc. Grenoble Workshop on Slow Neutrons* (1989).
39. G. Steigman, K. Olive, D.N. Schramm and M. Turner, *Phys. Lett. B* **176** (1986) 33.
40. ALEPH, DELPHI, L3, OPAL and MARK II collaboration papers, in *Proc. of the Int. High Energy Conference, Singapore* (1990).
41. D.N. Schramm and G. Steigman, *Phys. Lett. B* **141** (1984) 337.
42. E. Kolb and R. Scherrer *Phys. Rev. D* **25** (1982) 1481;  
E. Kolb, M.S. Turner, A. Chakravorty and D.N. Schramm, FERMILAB-PUB-91/28-A (1991) and *Phys. Rev. Lett.*, in press (1991).
43. S. Dimopoulos, R. Esmailzadeh, L. Hall and G. Starkman, *Astrophys. J.* **330**

- (1988) 545.
44. L. Brown and D.N. Schramm, *Ap.J.* **329** (1988) L103.
  45. D. Pilachowski, L. Hobbs and D. De Young, *Ap.J. Lett.* **345** (1989) L39.
  46. H. Reeves, K. Sato and M. Tarasawa, U. of Tokyo preprint (1989);  
Tarasawa, N. and Sato, K., U. of Tokyo preprint (1990).
  47. C. Alcock, G. Fuller, G. Mathews and B. Meyer, Livermore preprint (1990).
  48. F. Adams and K. Freese, MIT preprint (1990).
  49. L. Kawano, W. Fowler and R. Malaney, Caltech-Kellogg preprint (1990);  
J. Applegate, Columbia University preprint (1989).
  50. K. Freese and D.N. Schramm, *Nucl. Phys.* **B233** (1984) 167.
  51. M. Strauss, M. Davis and A. Yahil, U.C. Berkeley preprint (1989);  
N. Kaiser and A. Stebbins, CITA preprint (1990);  
E. Bertschinger, A. Dekel and A. Yahil, in *Proc. Blois Symposium on the Microwave Background* (1990);  
M. Rowan-Robinson and A. Yahil in *Proc. Rencontres de Moriond* (1989).
  52. A. Guth, *Phys. Rev. D* **23** (1981) 347;  
see also K. Olive, *Physics Reports* **190** (1990) 309-403 or A. Linde, *Particle Physics and Inflationary Cosmology*, (Harwood, New York, 1990).
  53. D. Hegyi and K. Olive, *K. Astrophys. J.* **303** (1986) 56.
  54. J. Gunn, talk at ITP Santa Barbara (1988);  
D. York, talk at University of Chicago (1988).
  55. G. Mathews and D.N. Schramm, *Astrophys. J.*, submitted (1991).
  56. L. Cowie, in *Proc. IUPAP Conference on Primordial Nucleosynthesis and Evolution of Early Universe, Tokyo, September 1990*, in press (1991).
  57. R. Bond and A. Szalay, *Proc. Texas Relativistic Astrophysical Symposium, Austin, Texas* (1982).
  58. M. Turner, F. Wilczek and A. Zee, *Phys. Lett. B* **125** (1983) 35 and **125** (1983) 519.
  59. M. Crawford and D.N. Schramm, *Nature* **298** (1982) 538.  
E. Witten, *Phys. Rev. D* **30** (1984) 272;  
see also C. Alcock and A. Olinto, *Ann. Rev. Nuc. Part. Phys.* **38** (1988) 161.
  60. Ya. Zeldovich, *Adv. Astron. and Astrophys.* **3** (1965) 241.
  61. B. Lee and S. Weinberg, *Phys. Rev. Lett.* **39** (1977) 165.
  62. H.-Y. Chiu, *Phys. Rev. Lett.* **17** (1966) 712.  
C.P. Hut, *Phys. Lett. B* **69** (1977) 85;  
K. Sato and H. Koyayashi, *Prog. Theor. Phys.* **58** (1977) 1775.
  63. S. Dimopoulos, R. Esmailzadeh, L. Hall and N. Tetradis, *Nucl. Phys. B*, submitted (1990).
  64. D. Brahm and L. Hall, *Phys. Rev. D* **41** (1990) 1067.
  65. K. Griest, M. Kamionkowski, and M. Turner, *Phys. Rev. D* **41** (1990) 3565.  
K. Olive and M. Srednicki, *M., Phys. Lett. B* **230** (1989) 78.
  66. D. Caldwell *et al.*, in *Proc. La Thuile*, ed. M. Greco;  
see also Kamiokande II Collaboration, in *Proc. IUPAP Symposium on Primordial Nucleosynthesis and Evolution of Early Universe, Tokyo, September 1990*, in press (1991).
  67. L. Krauss, *Phys. Rev. Lett.*, in press (1990);

- K. Griest and J. Silk, J., U.C. Berkeley preprint (1990).
68. J. Ellis, D. Nanopoulos, L. Roskowski and D.N. Schramm, *Phys. Lett. B* **245** (1990) 251-257.
  69. D. N. Schramm, in *Proc. 1988 Berkeley Workshop on Particle Astrophysics*, ed. E. Norman (World Scientific, Singapore, 1989).
  70. S.M. Faber and J.S. Gallagher, *Ann. Rev. Astron. and Astrophys.* **17** (1979) 135.
  71. E. Loh and E. Spillar, *Astrophys. J.* **329** (1988) 24.
  72. C. Hill, D.N. Schramm and J. Fry, *Comments Nucl. Part. Phys.* **19** (1989) 25.  
D.N. Schramm, in *Proc. IUPAP Symposium on Primordial Nucleosynthesis and Evolution of Early Universe, Tokyo, September 1990*, in press (1991).
  73. SAGE collaboration, in *Proc. Int. High Energy Physics Meeting, Singapore, August 1990*, in press (1991).
  74. G. Smoot, in *Proceedings of the Tokyo meeting on Primordial Nucleosynthesis, September 1990*. (1990).
  75. D. Schneider, M. Schmidt and J. Gunn, *Astron. J.* **98** (1989) 1951-1958.
  76. G. Efstathiou and M.J. Rees, *Mon. Not. Royal Astro. Soc.* **230** (1988) 5-11.
  77. K. Olive, *Physics Reports* **190** (1990) 309-403.
  78. E. Kolb and M. Turner, *The Early Universe*, (Addison Wesley, San Francisco, 1989).
  - 79 D.N. Schramm, in *Proc. of the 1990 Nobel Symposium, Graäftåvalen, Sweden* (1990).
  80. I. Wasserman, *Phys. Rev. Lett.* **57** (1986) 2234.
  81. W. Press, B. Ryden and D. Spergel, *Phys. Rev. Lett.* **64** (1990) 1084.
  82. C.T. Hill, D.N. Schramm, L.M. Widrow, in *Proc. of the XXVth Rencontres de Moriond at Les Arcs, Savoie, France, March 1990*, in press, and *Science*, submitted (1990).
  83. J. Frieman, C. Hill and R. Watkins, FERMILAB-PUB, in press (1991).
  84. M. Geller and J. Huchra, *Science* **246** (1989) 897.
  85. T. Broadhurst, R. Ellis, D. Koo and A. Szalay, *Nature* **343** (1990) 726.
  86. D.C. Koo and R.G. in *Towards Understanding Galaxies at Large Redshifts*, eds. F.G. Kron and A. Renzini (Kluwer Academic Publishers, Dordrecht, 1990) p. 209.
  87. A. Szalay, in: *Proc. IUPAP Conf. Primordial Nucleosynthesis and Evolution of Early Universe, Tokyo, September 1990* (1990).
  88. S. Faber and A. Dressler, in *Proc. Texas Symposium on Rel. Astrophysics, Dallas* (1988).
  89. N. Bahcall and R. Soneira, *Astrophys. J.* **270** (1983) 20-38.
  90. A. Klypin and A. Kopylov, *Soviet Astr. Lett.* **9** (1983) 41-46.
  91. J. Primack and A. Dekel, U. of California at Santa Cruz preprint (1990).
  92. A. Szalay and D.N. Schramm, *Nature* **314** (1985) 718-719.
  93. J. Huchra, Harvard University preprint (1990).
  94. S. Vandenburg and J. West, Dominion Observatory preprint (1990).
  95. G. Fuller and D.N. Schramm, University of California at San Diego preprint (1990).
  96. C. Hill, and G. Ross, *Phys. Lett. B* **203** (1988) 125.

97. S. Dimopoulos, private communication (1990).
98. S. P. Mikheyev and A. Yu Smirnov, *Nuovo Cimento* **9C** (1986) 17.
99. L. Wolfenstein, *Phys. Rev. D* **17** (1978) 2369.
100. J. Bahcall, *Neutrino Astrophysics* (Cambridge University Press, Cambridge, England and New York, 1989).
101. SAGE collaboration, Int. High Energy Meeting, Singapore, 1990 (1990).
102. M. Gell-Mann, P. Ramond and R. Slansky, R., in *Supergravity*, eds. P. Van Nieuwenhuizen and D. Freedman (North Holland Pub. Co., Dordrecht, 1979),
103. M. Dine, W. Fischler and M. Srednicki, *Phys. Lett. B* **104** (1981) 199.
104. W. Press, B. Ryden and D. Spergel, *Astrophys. J* **397** (1989) 590.
105. H. Hodges, *Phys. Rev. D* **37** (1988) 3052-5.
106. A.J. Stebbins and M.S. Turner, *Astrophys. J.* **339** (1989) L13.
107. L. Widrow, CfA/Harvard preprint (1990).
108. L. Kawano, *Phys. Rev. D* **41** (1990) 1013.
109. A. Massarotti, FERMILAB-PUB-90/77-A (1990) and *Phys. Rev. D*, in press (1991).
110. J. Ipser and P. Sikivie, *Phys. Rev. D* **30** (1984) 712.
111. G. Götz, FERMILAB-PUB-90/35-A (1990).
112. N. Turok, *Phys. Rev. Lett.* **63** (1989) 2625.
113. T.D. Lee, *Phys. Rev. D* **35** (1987) 3637.
114. V. Icke and R. van de Weygaert, *Quarterly Journal of the Royal Astronomical Society*, in press (1990).
115. P. Coles, *Mon. Notices Royal Astronomical Society*, submitted (1990).
116. E. Bertschinger, R. Scherrer, J. Vilenson, and A. Van Dalen, Ohio State preprint (1990).
117. M. Rees and D. Sciama, *Nature* **217** (1968) 511.
118. R.K. Sachs and A.M. Wolf, *Astrophys. J.* **147** (1967) 74.
119. M. Turner, R. Watkins and L. Widrow, L., *Astrophys. J. Lett.*, submitted (1990).
120. G. Götz and D. Nötzold, FERMILAB-PUB-89/235-A (1989).  
G. Götz and D. Nötzold, FERMILAB-PUB-89/236-A (1989).
121. K. Olive, D. Schramm, D. Thomas and T. Walker, U. of Minn. preprint (1991).
122. J.J. Simpson, *Phys. Lett. B* **174** (1986) 113;  
J.J. Simpson and A. Hime, *Phys. Rev. D* **39** (1989) 1825.
123. E. B. Norman, in *Proc. 14th Europhys. Conf. on Nucl. Phys., Bratislava, Czecko-Slovakia* (1990);  
B. Sur *et al.*, Lawrence Berkeley Laboratory preprint 30109 (1990).
124. A. Hime and N.A. Jelly, Oxford University preprint OUNP-91-01 (1990).
125. See, for example, discussion by F. Boehm, California Institute of Technology preprint (1991).
126. E.W. Kolb and M.S. Turner, FERMILAB-PUB-91/50-A (1991);  
L. Krauss, Yale preprint YCTP-P13-91 (1991);  
D. Caldwell and P. Langacker, UCSB-HEP 91-02 (1991);  
S.L. Glashow, Harvard University preprint HUTP-90/A075 (1990);  
S.K. Babu and R.N. Mohapatra, University of Maryland preprint UMO-PP-91-186 (1991);

- L. Dickson and Y. Nir, Stanford preprint SLAC-Pub-5496 (1991);  
 G. Ross, Oxford preprint (1991).  
 R. Barbieri and L. Hall, Lawrence Berkeley Laboratory preprint 30465 (1991);  
 T. Walker, Ohio State preprint (1991).
127. S. Dimopoulos, private communication (1991).
  128. K. Olive, D.N. Schramm and G. Steigman, *Nucl. Phys.* **180** (1981) 497-515.
  129. R. Gandhi and A. Burrows, *Phys. Lett. B* **242** (1990) 77.
  130. M.S. Turner, private communication (1991).
  131. R. Mayle, J. Wilson and D.N. Schramm, *Astrophys. J.* **318** (1987) 288.
  132. D.N. Schramm and J. W. Truran, *Phys. Rep.* **189** (1990) 89.
  133. For a summary of basic astrophysical constraints on neutrinos see E.W. Kolb, D.N. Schramm and M.S. Turner, in *Neutrino Physics*, ed. Klaus Winter, in press (Cambridge University Press, 1991).