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ENHANCED HEAVY-ELEMENT FORMATION IN BARYON-INHOMOGENEOUS BIG BANG MODELS

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ABSTRACT

We show that primordial nucleosynthesis in baryon-inhomogeneous big-bang models can lead to significant heavy-element production while still satisfying all of the light-element abundance constraints including the low lithium abundance observed in Population II stars. The parameters which admit this solution arise naturally from the process of neutrino-induced expansion of baryon inhomogeneities prior to the epoch of nucleosynthesis. These solutions entail a small fraction of baryons ($\lesssim 2\%$) in very high density regions with local baryon to photon ratio $\eta^h \simeq 10^{-4}$, while most baryons are at a baryon-to-photon ratio which optimizes the agreement with light-element abundances. This model would imply a unique signature of baryon inhomogeneities in the early universe, evidenced by the existence of primordial material containing heavy-element products of proton and alpha-burning reactions with an abundance of $[Z] \sim -6$ to -4.

Subject headings: early universe — nuclear reactions, nucleosynthesis, abundances

1. INTRODUCTION

Recently, there has been considerable interest in the possibility that primordial nucleosynthesis may have occurred in an environment in which the baryons were distributed inhomogeneously (Alcock, Fuller, & Mathews 1987; Applegate, Hogan, & Scherrer 1987, 1988; Fuller, Mathews, & Alcock 1988; Kurki-Suonio et al. 1988, 1990; Malaney & Fowler 1988; Boyd & Kajino 1989; Terasawa & Sato 1989a, b, c, 1990; Kajino & Boyd 1990; Kurki-Suonio & Matzner 1989, 1990; Mathews et al. 1990; Mathews, Schramm, & Meyer 1993; Kawano et al. 1991; Jedamzik, Fuller, & Mathews 1994; Thomas et al. 1994). Such studies were at first motivated by speculation (Witten 1984; Applegate & Hogan 1985) that a first-order quark-hadron phase transition at ($T \sim 100$ MeV) could produce baryon inhomogeneities as baryon number was trapped within bubbles of shrinking quark-gluon plasma. Even if this transition is not first order as some lattice-gauge theory calculations now seem to imply (e.g., Brown et al. 1990) such fluctuations might also have been produced by a number of other processes operating in the early universe such as late baryogenesis during the electroweak phase transition (Fuller et al. 1993), a TeV-scale $Z(3)$ QCD phase transition (Ignatius, Kajantie, & Rummukainen 1992), kaon condensation after the QCD transition (Nelson 1990), or magnetic fields from superconducting cosmic strings (Malaney & Butler 1989). Other mechanisms have also been proposed (see Malaney & Mathews 1993 for a recent review).

The first calculations (Alcock et al. 1987; Applegate et al. 1987; 1988; Malaney & Fowler 1988; Fuller et al. 1988) of the effects of such inhomogeneities were particularly interesting since they seemed to admit a much larger average universal baryon-to-photon ratio η than that allowed in the standard homogeneous big bang model (Wagoner, Fowler, & Hoyle 1967; Schramm & Wagoner 1977; Yang et al. 1984; Boesgaard & Steigman 1985; Olive et al. 1990; Walker et al. 1991; Smith, Kawano, & Malaney 1993) while still satisfying the light-element abundance constraints. Subsequent calculations (Kurki-Suonio et al. 1990; Malaney & Fowler 1988; Terasawa & Sato 1989a, b, c, 1990; Kajino & Boyd 1990; Kurki-Suonio & Matzner 1989, 1990; Mathews et al. 1990, 1993; Kawano et al. 1991; Jedamzik et al. 1994; Thomas et al. 1993) have shown, however, that when baryon diffusion is properly coupled with the dynamics of the system, the allowed values of η which are consistent with the light-element abundances are not significantly different than those of the standard homogeneous big bang.

Even with this constraint on η, however, an interest in baryon inhomogeneities has remained due to the possibility that such fluctuations might have produced an observable signature in the abundances of heavier elements such as beryllium and boron (Kajino & Boyd 1990; Malaney & Fowler 1989; Terasawa & Sato 1990; Kawano et al. 1991), intermediate mass elements (Kajino, Mathews, & Fuller 1990), or heavy elements (Applegate et al. 1988; Rauscher et al. 1993). Such possible signatures are also constrained, however, by the light-element abundances (e.g., Terasawa & Sato 1990). In particular, the Population II lithium abundance, Li/H $\sim 10^{-10}$, constrains the possible abundances of synthesized heavier nuclei to be quite small (e.g., Alcock et al. 1990; Terasawa & Sato 1990).

The purpose of this paper, however, is to point out a previously unexplored region of the parameter space for baryon-inhomogeneous big bang models in which substantial production of heavy elements is possible without violating any of the light-element abundance constraints. This solution

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arises naturally from one of the most likely scenarios for the production of baryon inhomogeneities in the first place. In fact, we discovered this solution not by looking for ways to optimize heavy-element production, but by exploring the natural evolution of large baryon number-density fluctuations. If such fluctuations are created in an early epoch, then they will likely contain a small fraction of the baryons. They will also be damped subsequently by neutrino-induced heating and expansion to a characteristic baryon-to-photon ratio $\sim 10^{-4}$ (Jedamzik & Fuller 1994; Jedamzik et al. 1994), providing the requisite environment for unique heavy-element nucleosynthesis.

2. BARYON INHOMOGENEITIES AND NEUTRINO INFLATION

Let us begin with a discussion of how baryon inhomogeneities might be generated in the early universe and how they are most likely to appear. The most likely scenario probably involves a first-order phase transition. It is not clear whether that transition is associated with the QCD epoch or an earlier time. Nevertheless, any process which couples to baryons in a spatially dependent way can produce baryon inhomogeneities (Malaney & Mathews 1993). For example, in the QCD transition, inhomogeneities would most likely be produced through a combination of the limited permeability of the phase boundary and the higher thermodynamic solubility of baryon number in the high-temperature phase (Witten 1984; Fuller et al. 1988; Kurki-Suonio 1988). The same might be true if inhomogeneities are produced by a kaon condensate (Nelson 1990) after the QCD transition.

If baryogenesis is associated with a first-order phase transition at the electroweak epoch (Kuzmin, Rubakov, & Shaposhnikov 1985; Cohen, Kaplan, & Nelson 1990; Dine et al. 1991; Turok & Zadrozny 1990; McLerran et al. 1991) then it may well be that the net baryon number is generated in a inhomogeneous fashion (Fuller et al. 1993). The presence of cosmic strings might also induce baryon inhomogeneities through electromagnetic (Malaney & Butler 1989) or gravitational interactions.

It is a common misconception that the most natural amplitude for baryon-number fluctuations is just that given by the thermodynamic ratio of equilibrium baryon densities in the two phases of a QCD transition. This as a number ~ 100 for the QCD transition (Witten 1984; Alcock et al. 1987). This value for the fluctuation amplitude, however, is unlikely. It would occur only if complete equilibrium were maintained in both phases until near the end of the phase transition followed by a sudden complete drop from equilibrium. This would require efficient mixing of baryon number in both phases and efficient transport of baryon number across the phase boundary until just near the end of the phase transition. Near the end of the phase transition, the transport of baryon number across the phase boundary cannot be efficient enough to establish chemical equilibrium. This is because the velocity of the phase boundary continuously increases. Efficient mixing of baryon number within the two phases will also be unlikely due to the short mean free path for baryons for $T \geq 50$ MeV. Furthermore if baryon mixing were this efficient throughout the phase transition, it would likely be sufficiently efficient after the phase transition to homogenize any baryon inhomogeneities long before the epoch of nucleosynthesis.

Mechanisms which take into account the efficiency for baryon transport, in particular, can lead to significant baryon inhomogeneities, and even in the most conservative conditions (e.g., Kurki-Suonio 1988) can produce a central baryon density far in excess of the thermodynamic equilibrium density (Fuller et al. 1988). If the baryon fluctuations are produced by such a mechanism, there is almost no limit as to how high the central baryon-to-photon ratio can become. However, for realistic estimates of baryon transport within the high-density phase and across the phase boundary, the total fraction of baryons remaining in high-density fluctuations tends to be small (Kurki-Suonio 1988; Fuller et al. 1988). The possibility of such a high baryon density in fact motivated Witten (1984) to suggest this mechanism as a means to form stable quark-matter agglomerates during the QCD phase transition if such matter were the ground state of baryonic matter at high density.

In any event, it is certainly possible that baryon inhomogeneities produced by any of the above mechanisms could reach a central baryon-to-photon ratio of $\eta \geq 10^{-4}$. This being the case, it is important to note the effect of neutrinos on any such large-amplitude baryon fluctuation (Hogan 1978, 1990; Heckler & Hogan 1993; Jedamzik & Fuller 1994). These overdense regions will have a higher pressure contribution from baryons than their surroundings. Hence, in order to maintain pressure equilibrium with the surrounding photon bath, any region of baryon overdensity will be at a slightly lower temperature than the baryon-depleted regions. However, since neutrinos have a longer mean free path than photons, they will penetrate into and heat the high-baryon-density regions causing them to expand and diminish the excess baryon density.

Several recent papers have explored this process in detail (Heckler & Hogan 1993; Jedamzik & Fuller 1994; Jedamzik et al. 1994). Jedamzik & Fuller (1994) found that this neutrino-induced damping of inhomogeneities is more pronounced as the density increases, such that any fluctuation with baryon-to-photon ratio in the high-density regions, $\eta^h \gtrsim 10^{-4}$, is damped to $\eta^h \approx 10^{-4}$. This is independent of its initial baryon density and holds true over a wide range of initial fluctuation radii. Fluctuations with η less than this are essentially unchanged.

This is illustrated in Figure 1 which shows the final baryon-to-photon ratio η_f after neutrino inflation as a function of initial comoving radius of a fluctuation. Lines are drawn corresponding to different initial baryon-to-photon ratio, η_i. The

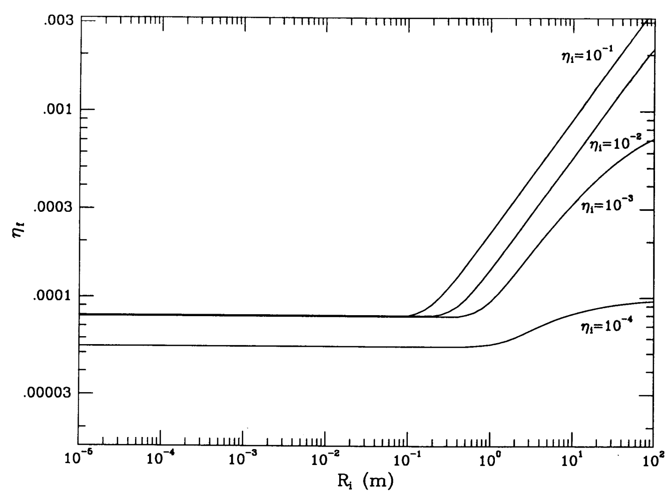

FIG. 1.—Final baryon-to-photon ratio, η_f, of neutrino-inflated fluctuations as a function of initial fluctuation radius (comoving at 100 MeV). Results are shown for four different initial values of the baryon-to-photon ratio, η_i.

degree to which the fluctuations are damped is related to the timescale during which the neutrino mean free path is of order the radius of the fluctuation. For larger fluctuation radii, this timescale becomes shorter due to the rapid increase in neutrino mean free path as the temperature decreases. Hence, if the initial radius of the fluctuation is too large, neutrinos cannot sufficiently damp the fluctuations to bring them to the limiting value of $\eta \sim 10^{-4}$.

The effects of neutrino inflation on very large amplitude baryon fluctuations with initial baryon-to-photon ratios of order unity ($\eta_i \sim 1$) are not shown in Figure 1. Fluctuations at such high baryon densities could correspond to stable or metastable low-entropy remnants like strange-quark-matter nuggets (Witten 1984). In this case heating by neutrinos at temperatures of $T \approx 50$ MeV will cause the evaporation of such nuggets rather than inflation. Only the very largest nuggets will survive complete evaporation (Alcock & Farhi 1985; Madsen, Heiselberg, & Riisager 1986). Evaporated nuggets can leave behind a fluctuation with a baryon-to-photon ratio significantly larger than $\eta \approx 10^{-4}$ even after neutrino inflation. The exact baryon distribution after evaporation will depend on the initial baryon number of the nuggets as well as on bulk properties of quark matter and strong interaction physics at the surface of the nuggets.

In any case, fluctuations with small initial radius and initial baryon-to-photon ratio less than unity will be inflated to a natural amplitude of $\eta^h \approx 10^{-4}$ by the time of primordial nucleosynthesis. The average baryon-to-photon value in the baryon-depleted regions η^l (or equivalently the total average baryon-to-photon ratio, η) and the total fraction of baryons in the high-density regions, f_b, can then be fixed by the model for forming the fluctuations or treated as free parameters to be fixed by the constraints from light-element nucleosynthesis. As noted above, models for the formation of inhomogeneities are consistent with small values of f_b. We will now show that a small value for f_b is also desired for agreement with light-element nucleosynthesis.

3. RESULTS

Fixing η at 10^{-4} in the high-density regions can avoid the overproduction of ^{7}Li. This was apparent already in the standard big bang calculations of Wagoner et al. (1967). This value for η is, however, much larger than has been considered in most standard or inhomogeneous big bang models since that early work. We can illustrate the effects of a high η^h value by first considering the simplest possible baryon inhomogeneous model, that in which the fluctuations are so widely separated that no significant baryon diffusion occurs before or during the epoch of primordial nucleosynthesis. In this limit the high-density and low-density regions can be treated (Wagoner et al. 1967; Wagoner 1973; Schramm & Wagoner 1977; Yang et al. 1984) as separate standard homogeneous big bang models with different local values for η. The nucleosynthesis yields from these regions can then be averaged after the epoch of nucleosynthesis. (We will also consider below the effects of baryon diffusion on small-scale fluctuations.)

The results of standard big bang nucleosynthesis (with three light neutrino species) at such high densities are summarized in Figure 2 where we have extended the big bang nuclear reaction network through mass $A = 28$. The point which has been overlooked in recent studies of baryon inhomogeneous models is that, for η greater than a few times 10^{-5}, ^{7}Li is destroyed by the $^7\mathrm{Li}(\alpha, \gamma)^{11}\mathrm{B}$ reaction. The ^{11}B thus formed is then rapidly con-

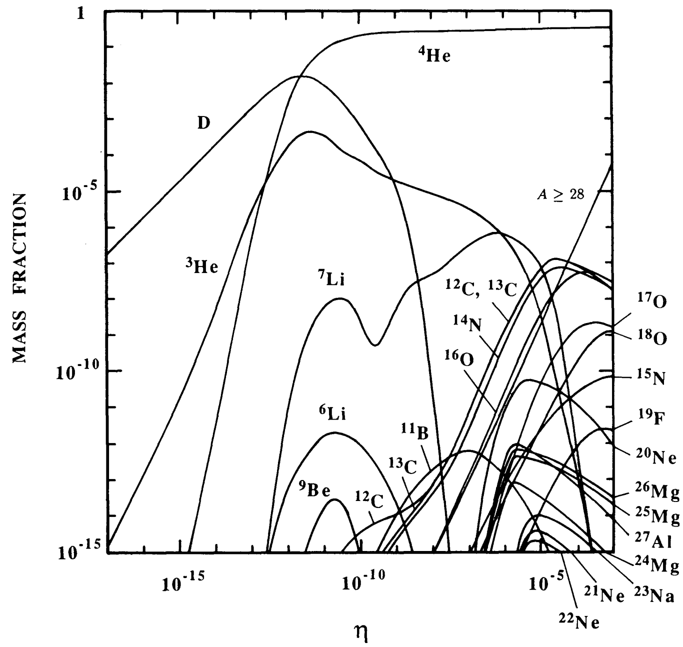

FIG. 2.—Nucleosynthesis yields as a function of baryon-to-photon ratio, η, for a reaction network extending from light nuclei to $A = 28$. These yields are for locally homogeneous regions sufficiently separated that baryon diffusion is insignificant. Note that for $\eta \gtrsim 10^{-4}$ the yields of nuclei heavier than helium are almost entirely in massive nuclei with essentially no production of other light elements.

sumed by a sequence of proton and alpha captures up to heavy nuclei reminiscent of an *rp*-process (Wallace & Woosley 1981). In fact, for $\eta \geq 10^{-4}$, almost all of the material with atomic mass heavier than ^{4}He is processed to nuclei beyond the end of our network at $A = 28$. Thus, we avoid the overproduction of ^{7}Li usually associated with averaging over η-values near the lithium minimum at $\eta \sim 2 \times 10^{-10}$ (e.g., Alcock et al. 1987).

It is, of course, still necessary to produce some deuterium, and in so doing, some ^{3}He and ^{7}Li will also be produced. There is also the problem that the ^{4}He abundance is significantly overproduced, that is, $Y_p \simeq 0.36$ for $\eta \simeq 10^{-4}$. The simultaneous solution to these constraints fixes allowed values of η^l and the fraction of baryons in the high-density region, $f_b = f_v(\eta^h/\eta)$, where f_v is the fraction of total volume occupied by the high-density regions and η is the average baryon-to-photon ratio.

We adopt constraints on observed light-element abundances from Walker et al. (1991) and Smith et al. (1993) such that in this simple model with no diffusion, the set of constraint equations become

$$0.22 \leq f_b Y_h + (1 - f_b) Y_l \leq 0.24 \, , \quad (1a)$$

$$\frac{\mathrm{D}}{\mathrm{H}} = \frac{X^h_{\mathrm{D}} f_b + (1 - f_b) X^l_{\mathrm{D}}}{2[X^h_{\mathrm{H}} f_b + (1 - f_b) X^l_{\mathrm{H}}]} \simeq (1 - f_b) \left(\frac{\mathrm{D}}{\mathrm{H}}\right)^l \geq 1.8 \times 10^{-5} \, , \quad (1b)$$

$$[\mathrm{D} + {}^3\mathrm{He}]/\mathrm{H} \simeq (1 - f_b)([\mathrm{D} + {}^3\mathrm{He}]/\mathrm{H})^l \leq 1.0 \times 10^{-4} \, , \quad (1c)$$

$$0.8 \times 10^{-10} \leq \mathrm{Li/H} \simeq (1 - f_b)(\mathrm{Li/H})^l \leq 2.3 \times 10^{-10} \, , \quad (1d)$$

where X_i is the mass fraction of species, i, and the superscripts h and l refer to quantities computed in the high and low baryon density regions, respectively. The approximate relations in

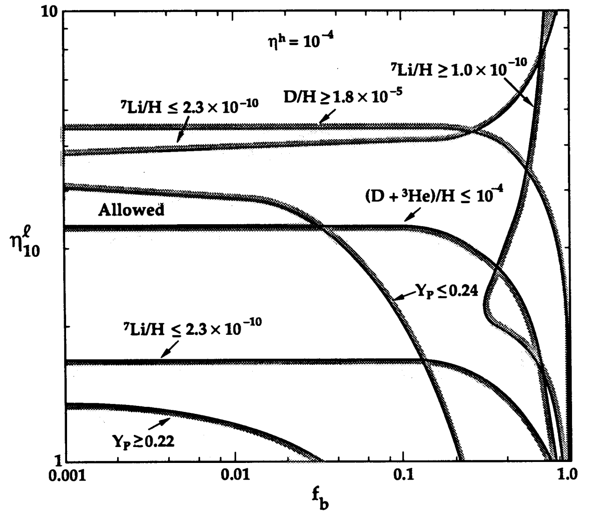

FIG. 3.—Light-element abundance constraints on the baryon-to-photon ratio in the low-density regions, η^l, as a function of the fraction, f_b, of baryons in high-density regions. Note that both quantities, η^l and f_b, are plotted on a logarithmic scale. For this model the high baryon density region is fixed by neutrino-induced inflation of the fluctuations to be $\eta^h = 10^{-4}$.

(1b)–(1d) are valid for small f_b and insignificant D, ^{3}He, or ^{7}Li production in the high baryon density regions.

Figure 3 shows allowed regions of the parameter space of η^l (in units of 10^{-10}) as a function of f_b. Clearly there is a region of f_b and η^l for which all of the light-element abundance constraints can be satisfied. This region is mostly defined by the upper limits to helium and [D + ^{3}He]/H. This figure does not include the extension of the allowed parameter space due to nuclear-reaction uncertainties, as is usually done in standard big bang studies (e.g., Walker et al. 1991; Smith et al. 1993). For example, if we had allowed for uncertainties in helium and deuterium production, the range of allowed η^l_{10} would expand to $2.8 \leq \eta^l_{10} \leq 4.0$ for small f_b, which is just the usually quoted standard big bang limit for three light neutrino species. The most stringent constraint on f_b is from the upper limit to the primordial helium abundance which necessitates that at most only ~2% of the baryons can be in high-density regions. If the uncertainty in the neutron half-life is also included, this upper limit increases to as high as $f_b \lesssim 0.1$.

It is easy to see from an inspection of Figure 2 why this region of the parameter space is allowed. By placing almost all of the baryons (≳98%) in a low-density region for which η^l is close to the optimum standard big bang value, a good fit to all of the light-element abundances is guaranteed.

So we see that in this large-separation limit for inhomogeneous models, all of the light-element constraints can be satisfied. Even though this constraint requires that only a small fraction of the baryons be in high-density regions, we still find a potentially interesting (though not necessarily practically observable) signature from inhomogeneous nucleosynthesis. For $\eta^h \simeq 10^{-4}$ and $f_b \sim 0.01$, the mass fraction Z of heavy ($A \geq 28$) or CNO nuclei is $Z \sim 3 \times 10^{-9}$ ($[Z] \simeq -6.5$) after averaging with the low-density regions.

The mass fraction in nuclei beyond the end of our network at $A = 28$ could reach as high as $[Z] \sim -4$ for $\eta^h \sim 10^{-3}$ and $f_b \sim 0.01$. A baryon-to-photon ratio as high as $\eta^h \sim 10^{-3}$ after the epoch of neutrino inflation requires an initially large-amplitude fluctuation of large radius, or equivalently a large baryonic mass $M_b \sim 5 \times 10^{-11}$ $M_\odot$ (see Fig. 1). Another possibility to obtain a local baryon-to-photon ratio during primordial nucleosynthesis in excess of $\eta^h \sim 10^{-4}$ is from the evaporation of low-entropy nuggets. The effects of strange-quark-matter nuggets on primordial nucleosynthesis have been considered (Schaeffer, Delbourgo-Salvador, & Audouze 1985; Madsen & Riisager 1985), but the evaporation of nuggets before primordial nucleosynthesis and the possibility of significant heavy-element production during primordial nucleosynthesis has not been addressed before.

It is interesting to note that a small fraction of baryons initially within large-mass fluctuations or low-entropy nuggets might produce a mass fraction in heavy elements during primordial nucleosynthesis, in excess of $[Z] \sim -4$. Thus observation (e.g., Beers, Preston, & Shectman 1992) of very low metallicity objects ($[Z] \lesssim -4$) can potentially constrain the existence of large-mass fluctuations or low-entropy remnants in the early universe. The existence of a floor in primordial heavy-element abundances could even provide a nucleosynthesis signature of the physical properties of the fluctuations.

In any case, fluctuations of moderate amplitude and size will be damped by neutrino inflation to a baryon-to-photon ratio of $\eta^h \sim 10^{-4}$. Hence, $[Z] \sim -6$ is probably a good characteristic abundance for the sum of heavy-element abundances produced by inhomogeneous nucleosynthesis.

In this case a possible abundance diagnostic signature includes C/O and N/O ratios (Kajino et al. 1990). A high C/O or N/O ratio is contrary to what one might expect from a first generation of massive stars, but does appear for big bang models with $10^{-5} \lesssim \eta^h \lesssim 10^{-4}$. For $\eta^h \gtrsim 10^{-4}$, one should see an anomalously low [O/Fe] ratio, assuming that the $A \geq 28$ nuclei are burned to Fe. We realize, of course, that detecting abundances this low is extremely difficult and probably not possible with existing techniques. Nevertheless, it may eventually be possible to measure accurately abundances this low, in which case the presence or absence of baryon inhomogeneities might be explorable.

3.1. *Effects of Baryon Diffusion*

The results shown in Figure 3, of course, correspond to only one possible model in which the separation between fluctuations is so large that baryon diffusion is insignificant. We have also explored models in which neutron diffusion has been included but for a reaction network limited to $A \leq 12$. We have considered models in which the proper separation distance (at $T = 100$ MeV) was varied between the QCD horizon (10^4 m) and a small fraction thereof (0.1 m) for fluctuations with fixed η^h and η^l. Alternatively, we have considered effects of baryon diffusion as a function of baryon mass in the fluctuation, M_b, and fraction f_b of the baryons in the high-density regions for fixed η^l.

Conventionally a regular lattice of fluctuations is described by four variables, the mean separation between adjacent fluctuations, r, as well as three more variables. These can be chosen as the volume fraction in the high-density region, f_v, the ratio of densities in high-density region and low-density region, R, and the average baryon-to-photon ratio, η (Fuller et al. 1988). In the present analysis of high-amplitude fluctuations, which include only a small fraction of the total baryons ($f_b \ll 1$), it is convenient to characterize the fluctuation by baryon mass instead of the mean separation distance. Varying f_b at fixed M_b is equivalent to varying the separation distance. Varying M_b at fixed f_b corresponds to a self-similar variation of fluctuation radius and separation distance.

A conversion between the mean separation and baryon mass is easily obtained, from

$$r \approx 10 \text{ m}\left(\frac{M_b}{10^{-21} M_\odot}\right)^{1/3}\left(\frac{\eta^l}{2.8 \times 10^{-10}}\right)^{-1/3}\left(\frac{f_b}{0.01}\right)^{-1/3}, \quad (2)$$

where r is the proper mean separation between fluctuations (at $T = 100$ MeV), and M_b is the baryon mass within a high-density fluctuation. In equation (2) and what follows it is straightforward to convert from proper length at $T_0 = 100$ MeV to proper length at temperature T. To a good approximation, one can do this by simply assuming that the scale factor $a \propto T_0/T$ for all temperatures $T > 0.1$ MeV. The baryon mass is related to the initial (final) fluctuation baryon-to-photon ratio, η_i (η_f), and fluctuation radius, R_i (R_f), as follows:

$$M_b \approx 4 \times 10^{-13} M_\odot \eta_i^h R_i^3 = 4 \times 10^{-13} M_\odot \eta_f^h R_f^3, \quad (3)$$

where the proper ($T = 100$ MeV) fluctuation radius R is in meters. Note that the product of baryon-to-photon ratio times the cube of the fluctuation radius is not changed by neutrino inflation because baryon number within the fluctuation is conserved.

With the help of equation (3) and Figure 1 the reader is able to determine the mass of a fluctuation, M_b, its baryon-to-photon ratio after neutrino inflation, η_f, and radius, R_f, assuming η_i and R_i.

In order to identify the mass scales for which baryon diffusion affects fluctuations before and during primordial nucleosynthesis, the baryon diffusion length, d, should be compared to the radius of the fluctuation after neutrino inflation, R_f. Baryon diffusion in high-density regions is limited by neutron-proton scattering (Applegate et al. 1987; Jedamzik & Fuller 1994). The proper (referred to $T_0 = 100$ MeV) baryon diffusion length at temperature T, that is, the average distance a neutron diffuses between high temperatures and temperature T, is given by

$$d(T) \approx 3 \times 10^{-2} \text{ m}\left(\frac{\eta_f^h}{10^{-4}}\right)^{-1/2}\left(\frac{T}{500 \text{ keV}}\right)^{-5/4}. \quad (4)$$

Clearly, the baryon diffusion length is smaller for a larger value of the baryon-to-photon ratio in the fluctuation. For a fluctuation of baryon mass, $M_b \lesssim 10^{-21} M_\odot$, the baryon diffusion length at $T \sim 500$ keV is comparable to the fluctuation size, and the fluctuations, therefore, vanish by diffusion before nucleosynthesis begins.

For mass scales $M_b > 10^{-21} M_\odot$, for which the fluctuations can survive to the epoch of nucleosynthesis, the abundance yields from models which include baryon diffusion are not that much different from Figure 2. This is because in high-density regions with $\eta^h \sim 10^{-4}$, nuclear reactions can begin to synthesize light elements at temperatures as high as $T \approx 300$–400 keV. Thus, in the high-density regions free neutrons are incorporated into helium soon after the freeze-out of weak reactions. As the weak reactions freeze out, neutrons and protons are no longer rapidly converted from one into another. Hence, they obtain separate identities.

Because the neutrons are absorbed so quickly after weak freeze-out, there is not much time for neutron diffusion to produce large variations in the neutron-to-proton ratio. Even though some free neutrons in the low-density regions can diffuse back into the high-density regions (Malaney & Fowler 1988), this back diffusion does not significantly affect the light-element abundances when $f_b \ll 1$. This is because light elements are produced in low-density regions. This influx of neutrons into the high-density regions may somewhat affect the yields of heavy elements which are produced there. For the most part, however, the neutron diffusion length in the high baryon density regions is short, and the final abundance yields should not differ much from that obtained by simply averaging over regions with different η and fixed n/p ratio as in Figure 2.

The results of numerical studies including the effects of baryon diffusion are presented in Figures 4 and 5. These figures illustrate that fluctuations more widely separated than a proper distance of $r = 100$ m (at $T = 100$ MeV) or of a mass scale $\geq 10^{-18} M_\odot$ produce nucleosynthesis yields which are indistinguishable from those of Figure 3 (Mathews et al. 1990). Similarly, models in which the proper separation distance (at $T = 100$ MeV) is $r \lesssim 5m$ ($M_b \leq 10^{-21} M_\odot$) are indistinguishable from a homogeneous standard big bang model with the same average η. The reason is that baryon diffusion prior to nucleosynthesis damps out the inhomogeneities (Mathews et al. 1990). The maximum effects on the nucleosynthesis yields occur for proper separation distances (at $T = 100$ MeV) between 10 and 60 m where the separation roughly equals the neutron diffusion length during primordial nucleosynthesis

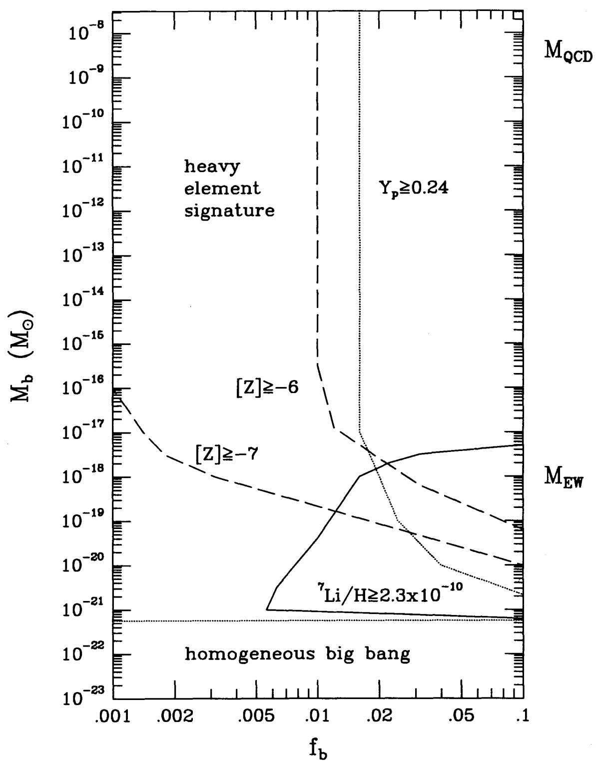

FIG. 4.—Heavy-element nucleosynthesis signatures of high-amplitude fluctuations ($\eta_i \geq 10^{-4}$, $\eta^l = 2.8 \times 10^{-10}$). Identified are regions of the total baryon mass M_b (in units of $M_\odot$) contained in a single fluctuation vs. the fraction of baryons in these high-density fluctuations, f_b. Also identified on this figure are the total baryon masses within the horizon for the electroweak transition (M_{EW}) and the QCD transition (M_{QCD}). For $M_b \gtrsim 10^{-18}$ (roughly the electroweak baryon mass scale) and $f_b \lesssim 0.02$, heavy elements are formed without overproducing lithium or helium. For less massive fluctuations, $M_b \lesssim 10^{-21}$, it is also possible to satisfy the lithium and helium constraints, but in this case baryon diffusion erases the inhomogeneities before the onset of primordial nucleosynthesis so that the results are indistinguishable from the homogeneous standard big bang and no heavy-element signature emerges.

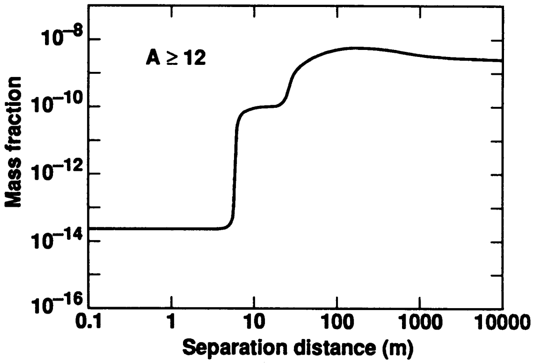

FIG. 5.—Example of the effects of baryon diffusion on heavy-element abundances ($A \geq 12$). Shown are average yields for heavy elements as a function of proper separation distance in m at $T = 100$ MeV. For this figure we have arbitrarily set $f_b = 0.0001$, $\eta^l = 2.8 \times 10^{-10}$, and $\eta^h = 1.3 \times 10^{-5}$. At separations less than 6 m, the results are indistinguishable from a standard homogeneous big bang with $\eta \sim 2.8 \times 10^{-10}$ because baryon diffusion before nucleosynthesis homogenizes the fluctuation. For separation distances greater than ~ 100 m, the results are essentially identical to Fig. 2. A transition between the two limits occurs when the neutron diffusion length during nucleosynthesis is comparable to the separation distance.

(Applegate et al. 1988; Mathews et al. 1990; Terasawa & Sato 1990).

Figure 4 identifies the region in the M_b (total baryon mass of the fluctuation) versus f_b plane in which heavy-element nucleosynthesis may imply a signature of high-amplitude fluctuations. To produce this figure we used the inhomogeneous nucleosynthesis code of Jedamzik et al. (1994) which includes all relevant hydrodynamic and radiation transport effects. We evolved baryon-number fluctuations between $T = 100$ MeV and the end of primordial nucleosynthesis at $T \simeq 5$ keV. The effects of neutrino inflation at high temperatures are included in the numerical treatment.

The total baryon mass within the horizon for the electroweak transition ($M_{\rm EW}$) and the QCD transition ($M_{\rm QCD}$) is also identified in Figure 4. For $M_b \gtrsim 10^{-18}$ (of order the electroweak baryon mass scale) and $f_b \lesssim 0.02$ heavy elements are formed without overproducing lithium or helium. For less massive fluctuations, $M_b \sim 10^{-21}$, it is also possible to satisfy the lithium and helium constraints, but in this case baryon diffusion erases the inhomogeneities before the onset of primordial nucleosynthesis so that the results are indistinguishable from the homogeneous case.

The effects of diffusion on the yields of light element abundances is not very large even at optimum separation distances or masses. This is because the high-density regions are constrained to contain only a small fraction of the total baryons. The main contribution to the light-element abundances is from the low-density regions which are not much affected by diffusion from the high-density regions.

The effects of baryon diffusion on the various light-element abundances are briefly summarized as follows:

Y_p—Independent of η^h, the helium mass fraction decreases by $\Delta Y_p \simeq 0.0005$ as the separation is decreased from 100 to 6 m. (For the remainder of this section all quoted length scales are proper lengths measured at $T = 100$ MeV.) As the separation is further decreased from 6 to 3 m, the helium mass fraction decreases by $\Delta Y_p \simeq 0.002$. This change is small compared to the uncertainty in the upper limit to Y_p.

D/H, (D + ^{3}He)/H—Because these light elements are made almost exclusively in the low-density regions which are largely uneffected by baryon diffusion, their only dependence upon the diffusion length is to decrease by about 4% as the diffusion length is varied between 50 and 5 m. Again, this change is insignificant compared to the uncertainties in the abundance constraints themselves.

^{7}Li—Lithium is the light element most significantly affected by neutron diffusion. This is largely because of ^{7}Be production at the fluctuation boundary (Mathews et al. 1990). This causes the lithium abundance to increase to a maximum of ^{7}Li/H $\sim 2 \times 10^{-10}$ for a separation of 10 m. For distances less than 6 m or greater than 100 m, the ^{7}Li abundance remains at the value on Figure 1 corresponding to whatever η^l-value is chosen.

^{9}Be, ^{11}B—Since there has been some discussion of the production of these light nuclei in recent literature, a mention of their production in the present models should be made here. There can be no production of Be in these models, and almost no ^{11}B production [$X(^{11}\rm B) \lesssim 10^{-18}$] even with an optimized separation distance (which, like lithium, can only increase boron by less than a factor of 2). Therefore, the abundances of these elements can not be a diagnostic for the kind of inhomogeneities considered here.

$A \geq 12$—By far the yield which is most sensitive to the diffusion efficiency is that of the heavy-element abundances. This sensitivity is illustrated in Figure 5. For separation distances less than $\simeq 6$ m the yields are no different than a standard homogeneous big bang with the same average value of η. In this case most of the yield is in ^{12}C at an insignificant abundance relative to hydrogen of less than 10^{-14} (see Fig. 2). However, as the diffusion length is increased to $\simeq 100$ m, the heavy-element abundance quickly increases by five orders of magnitude, due largely to alpha and proton captures in the high-density regions. For $\eta^h \gtrsim 10^{-4}$, heavy elements are mostly produced in the form of iron and an anomalously low [O/Fe] would be a diagnostic abundance signature. For $10^{-5} \lesssim \eta^h \lesssim 10^{-4}$, anomalously high [C/O] and [N/O] ratios are expected (Kajino et al. 1990). Any observation suggesting the presence of primordial heavy elements with abundances as large as $[Z] \gtrsim -6$ could be a strong verification of the presence of baryon inhomogeneities in the early universe.

4. CONCLUSIONS

We have shown that there exists a plausible region of the parameter space for baryon-inhomogeneous primordial nucleosynthesis models in which all of the light-element abundance constraints are satisfied, yet which also leads to the unique production of potentially observable abundances of heavy nuclei. Although the abundances of these heavy nuclei are small, they can be as much as six to eight orders of magnitude larger than in the homogeneous big bang. The region of the parameter space in which heavy-element abundances are generated is naturally motivated by models for the formation of baryon inhomogeneities, for example, during a first-order phase transition, and by numerical simulations of the effect of neutrino-induced heating and expansion of baryon inhomogeneities prior to the epoch of primordial nucleosynthesis. Hence, this region of the parameter space should be taken seriously.

The conditions most likely to result from the expansion of a high baryon density region by neutrino inflation, which also satisfy the light-element abundance constraints, imply that a mass fraction as large as $[Z] \sim -4$ to -6 could be in the form of nuclei with $A > 28$ and/or CNO nuclei after primordial nucleosynthesis. Thus, a search for a lower limit to the abundance ratios of CNO or heavy nuclei to the light primordial abundances could some day be a definitive indicator of the presence or absence of large-amplitude baryon inhomogeneities in the early universe.

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REFERENCES

Alcock, C. R., Dearborn, D. S., Fuller, G. M., Mathews, G. J., & Meyer, B. S. 1990, Phys. Rev. Lett., 64, 2607
Alcock, C. R., & Farhi, E. 1985, Phys. Rev. D, 32, 1273
Alcock, C. R., Fuller, G. M., & Mathews, G. J. 1987, ApJ, 320, 439
Applegate, J. H., & Hogan, C. J. 1985, Phys. Rev. D, 31, 3037
Applegate, J. H., Hogan, C. J., & Scherrer, R. J. 1987, Phys. Rev. D, 35, 1151
———. 1988, ApJ, 329, 572
Beers, T. C., Preston, G. W., & Shectman, S. A. 1992, AJ, 103, 1987
Boesgaard, A. M., & Steigman, G. 1985, ARA&A, 23, 319
Boyd, R. N., & Kajino, T. 1989, ApJ, 336, L55
Brown, F. R., Butler, F. P., Hong Chen, & Christ, N. H. 1990, Phys. Rev. Lett., 65, 2491
Cohen, A. G., Kaplan, D. B., & Nelson, A. E. 1990, Phys. Lett. B, 245, 561
Dine, M., Huet, P., Singleton, R., & Susskind, L. 1991, Phys. Lett. B, 257, 351
Fuller, G. M., Mathews, G. J., & Alcock, C. R. 1988, Phys. Rev. D, 37, 1380
Fuller, G. M., Jedamzik, K., Mathews, G. J., & Olinto, A. 1993 Phys. Lett., submitted
Heckler, A., & Hogan, C. J. 1993 Phys. Rev. D, 47, 4256
Hogan, C. 1978, MNRAS, 185, 889
———. 1990, in Primordial Nucleosynthesis, ed. W. J. Thompson, B. W. Carney, & H. J. Karwowski (Singapore: World Scientific), 15
Ignatius, J., Kajantie, K., & Rummukainen, K. 1992, Phys. Rev. Lett., 68, 737
Jedamzik, K., & Fuller, G. M. 1994, ApJ, in press
Jedamzik, K., Fuller, G. M., & Mathews, G. J. 1994, ApJ, in press
Kajino, T., & Boyd, R. N. 1990, ApJ, 359, 267
Kajino, T., Mathews, G. J., & Fuller, G. M. 1990, ApJ, 364, 7
Kawano, L. H., Fowler, W. A., Kavanagh, R. W., & Malaney, R. A. 1991, ApJ, 372, 1
Kurki-Suonio, H. 1988, Phys. Rev. D, 37, 2104
Kurki-Suonio, H., & Matzner, R. A. 1989, Phys. Rev. D, 39, 1046
———. 1990, Phys. Rev. D, 42, 1047
Kurki-Suonio, H., Matzner, R. A., Centrella, J., Rothman, T., & Wilson, J. R. 1988, Phys. Rev. D, 38, 1091
Kurki-Suonio, H., Matzner, R. A., Olive, K. A., & Schramm, D. N. 1990, ApJ, 353, 406
Kuzmin, V. A., Rubakov, V. A., & Shaposhnikov, M. E. 1985, Phys. Lett. B, 155, 36
Madsen, J., Heiselberg, H., & Riisager, K. 1986, Phys. Rev. D, 34, 2947
Madsen, J., & Riisager, K. 1985, Phys. Lett. B, 158, 208
Malaney, R. A., & Butler, M. N. 1989, Phys. Rev. Lett., 62, 117
Malaney, R. A., & Fowler, W. A. 1988, ApJ, 333, 14
———. 1989, ApJ, 345, L5
Malaney, R. A., & Mathews, G. J. 1993, Phys. Rep., 229, 145
Mathews, G. J., Meyer, B. S., Alcock, C. R., & Fuller, G. M. 1990, ApJ, 358, 36
Mathews, G. J., Schramm, D. N., & Meyer, B. S. 1993, ApJ, 404, 476
McLerran, L., Shaposhnikov, N., Turok, N., & Voloshin, M. 1991, Phys. Lett. B, 256, 451
Nelson, A. 1990, Phys. Lett. B, 240, 179
Olive, K. A., Schramm, D. N., Steigman, G., & Walker, T. P. 1990, Phys. Lett. B, 236, 454
Rauscher, T., Applegate, J. H., Cowan, J. J., Thielemann, F. K., & Wiescher, M. 1993, ApJ, submitted
Schaeffer, R., Delbourgo-Salvador, P., & Audouze, J. 1985, Nature, 317, 407
Schramm, D. N., & Wagoner, R. V. 1977, Ann. Rev. Nucl. Part. Sci., 27, 37
Smith, M. S., Kawano, L. H., & Malaney, R. A. 1993, ApJS, 85, 219
Terasawa, N., & Sato, K. 1989a, Prog. Theor. Phys., 81, 254
———. 1989b, Phys. Rev. D, 39, 2893
———. 1989c, Prog. Theor. Phys., 81, 1085
———. 1990, ApJ, 362, L47
Thomas, D., Schramm, D. N., Olive, K. A., Mathews, G. J., Meyer, B. S., & Fields, B. 1994, ApJ, submitted
Turok, N., & Zadrozny, J. 1990, Phys. Rev. Lett., 65, 2331
Wagoner, R. V. 1973, ApJ, 179, 343
Wagoner, R. V., Fowler, W. A., & Hoyle, F. 1967, ApJ, 148, 3
Walker, T. P., Steigman, G., Schramm, D. N., Olive, K. A., & Kang, H. 1991, ApJ, 376, 51
Wallace, R. K., & Woosley, S. E. 1981, ApJS, 45, 389
Witten, E. 1984, Phys. Rev. D, 30, 272
Yang, J., Turner, M. S., Steigman, G., Schramm, D. N., & Olive, K. A. 1984, ApJ, 281, 493