On the origin of the light elements ($Z < 6$)

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The author reviews the status of our understanding of nucleosynthesis of the light nuclei ($Z < 6$). The standard view today is that these elements are, for the most part, generated by two different processes: first, thermonuclear reactions in the early universe (big-bang nucleosynthesis or (BBN), and second, galactic cosmic-ray-induced spallation reactions (GCR) in cold interstellar atoms. A third contribution comes from stellar processes. The arguments in favor of this view are presented. Numerous astrophysical and cosmological implications are discussed, such as the baryonic density, the possible existence of baryonic dark matter and of nonbaryonic (exotic) matter, the constraints imposed on new particle physics, the leptonic number of the universe, the increase in cosmic entropy since primordial nucleosynthesis, and the constancy of the "constants" of physics.

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I. INTRODUCTION

Shortly after World War II, George Gamow (1946) and his collaborators (Apher *et al.*, 1948) considered the possibility that all chemical elements might be generated by a long chain of nucleon captures in the cooling primordial universe. The absence of stable nuclei with mass 5 and 8 constituted a fatal flaw of this scheme. Soon afterwards, Fred Hoyle and his collaborators (Burbidge *et al.*, 1957) proposed the idea of a stellar origin for the chemical elements. While this idea proved to be correct for the heavy nuclei, from carbon to uranium, it ran into difficulty in accounting for the lighter nuclei, given their short lifetimes in stellar interiors. Other processes were clearly needed to account for the low-Z element abundances.

One requirement for the selection of the appropriate formation mechanism is that these nuclei not remain at high temperature for any length of time. Primordial nucleosynthesis (BBN) generates D, ^{3}He, ^{4}He, and ^{7}Li. Cosmic matter's rapid rate of cooling guarantees their survival. GCR spallation creates ^{6}Li, ^{7}Li, ^{9}Be, ^{10}B, and ^{11}B directly in cold interstellar space. Stellar winds and explosions are the mechanisms by which stellar-produced ^{3}He, ^{4}He, and ^{7}Li are ejected into space.

II. NUCLEAR PHYSICS OVERVIEW

Many features of the universal abundance curve for chemical elements can be qualitatively understood through a knowledge of their nuclear properties. The iron peak (at $A=56$) corresponds to the most stable nuclear configuration. The secondary peaks correspond to nuclei with magic numbers of neutrons (at $N=50$, 82, and 128) and also to the light nuclei with an integer number of alpha particles (at $A=12$, 16, 20, 24, and 28). In this respect it is most informative to begin our study of the origin of the light elements with a brief review of their nuclear properties.

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Two important factors play a crucial role in the nucleosynthesis of the light nuclei: their small electric charge ($Z<6$) and the large binding energy of alpha clusters (2 protons and 2 neutrons) relative to other nuclear arrangements.

The Coulomb repulsion factor, at thermal energies much smaller than the Coulomb barrier, governs the destruction rate of the light nuclei by proton capture at stellar temperatures. The observed cross sections are given in Fig. 1 as a function of proton energy. Integrated over thermal distributions, they yield the threshold temperature above which a given nucleus cannot survive typical stellar conditions. Because of its low Z and low nuclear stability, deuterium is doubly fragile. It disappears at ≈ 0.5 million degrees. Next come ^{6}Li at ≈ 2.0; ^{7}Li at ≈ 2.5; ^{9}Be at ≈ 3.5; ^{11}B at ≈ 5.0; and ^{10}B at ≈ 5.3 (all in millions of degrees). As a result, the light elements (except ^{4}He and, to a certain extent, ^{3}He) do not resist the heat of stellar interiors. In view of the particle instability of ^{4}Li and ^{5}Li, the helium isotopes cannot be destroyed in proton-induced reactions, but only by helium-induced reactions, resulting in an appreciably higher Coulomb barrier energy.

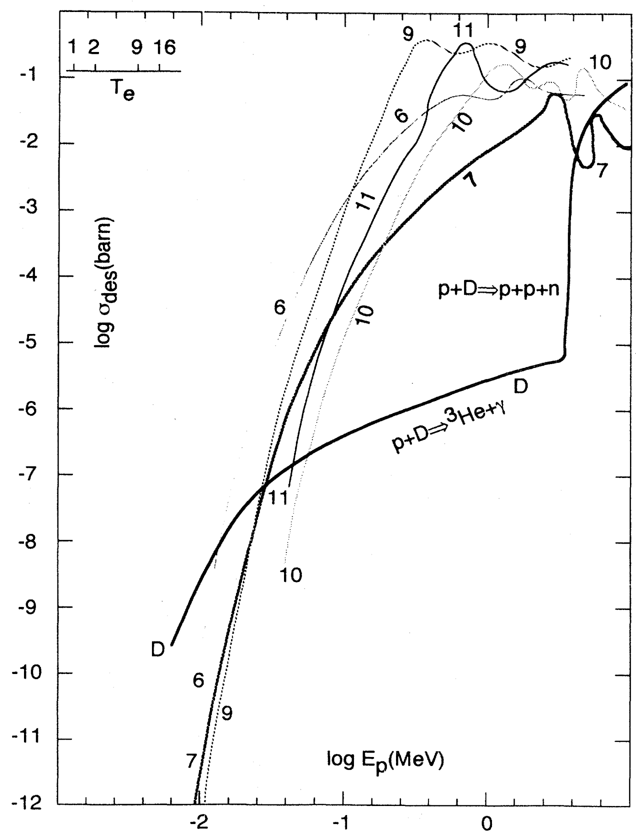

FIG. 1. Experimental cross sections for the destruction of D, ^{6}Li, ^{7}Li, ^{9}Be, ^{10}B, and ^{11}B as a function of incident-proton energies. Note the very small scale of the coordinate. The inset in the upper left shows the temperature in million degrees, for which the Gamow energies of Be are those of the abscissa. For example, at 10^7 K the Gamow energy of Be (or Li or B) is ≈ 10 keV.

The second important nuclear property is the large binding energy of ^{4}He, due to the large pairing effect of nuclear forces when the nucleons are paired four by four: neutron-proton; spin up, spin down. As a result, every nucleus in this mass range has a rather precarious stability toward a rearrangement involving ^{4}He nuclei. No mass-5 nucleus manages to be stable; the lifetimes are $\approx 10^{-21}$ sec. The isotopes ^{8}Li and ^{8}Be are beta unstable with respect to ^{8}Be, which quickly ($\approx 10^{-16}$ sec) breaks into two alphas.

The nuclear stability situation in the mass-9 range is deeply marked by the alpha stability. ^{9}B is unstable against $2\alpha+p$ ($\approx 10^{-19}$ sec), but ^{9}Be barely escapes disintegration (very small binding energy). This weak stability is reflected in the fact that the endothermic Q values corresponding to the formation of ^{9}Be by spallation reactions are remarkably large and that the exothermic Q values corresponding to its destruction are small. It is also reflected in the fact that it has only one "bound" state; all the excited states are unstable against particle breakup. These facts are instrumental in explaining its low spallation cross section and hence its low natural abundance (one of the lowest in nature). The isotopes ^{6}Li, ^{7}Li and ^{10}B, ^{11}B fare better, but all remain comparatively fragile; in high-energy proton-induced reactions, they all quickly break into residues involving alpha particles.

The high binding energy of the ^{4}He is also responsible for the fact that, at all but the lowest temperature, hydrogen is transformed all the way to ^{4}He. The nuclei D and ^{3}He are intermediate steps in this chain of reactions; comparatively small amounts remain at the end of the process. These facts will dominate the scenario of big-bang nucleosynthesis and also of main-sequence stellar energy generation.

The influence of these nuclear properties on the formation rate of the light elements is reflected in their relative

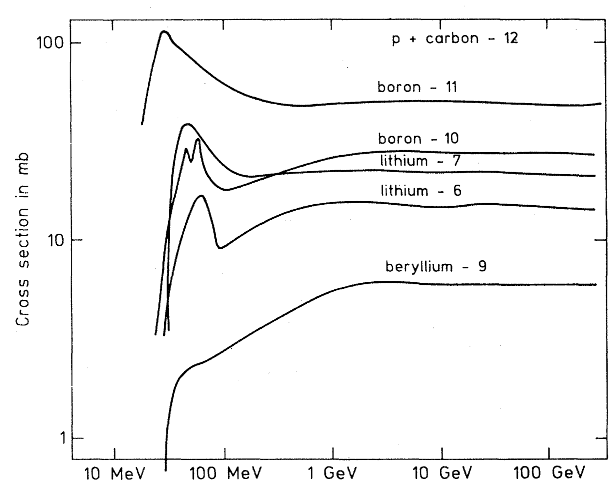

FIG. 2. Spallation cross sections of protons on ^{16}O as a function of energy. Note that, from a few hundred MeV upward, the cross sections are almost constant. Note also the decreasing sequence of B, Li, and Be.

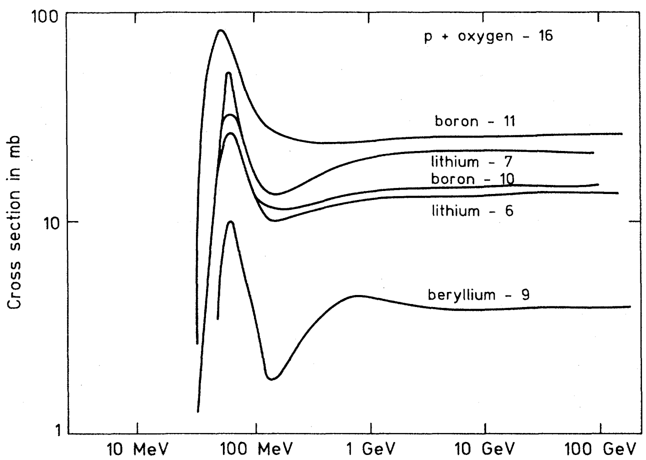

FIG. 3. Spallation cross sections of protons on ^{12}C as a function of energy. Note that, from a few hundred MeV upward, the cross sections are almost constant. Note also the decreasing sequence of B, Li, and Be.

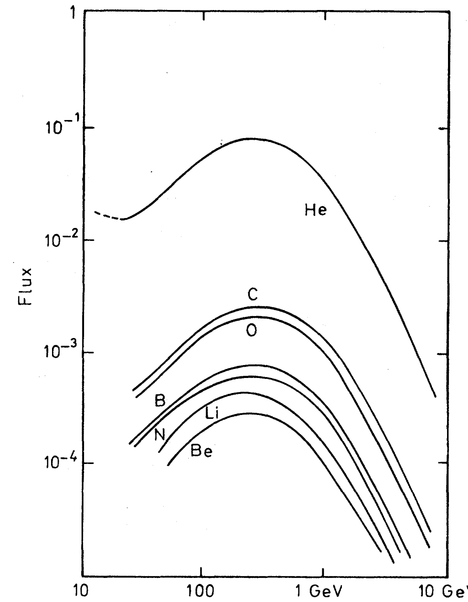

FIG. 5. Galactic cosmic-ray fluxes He, O, C, N, B, Li, and Be as a function of energy per nucleon. Note that B is more abundant than N. B is higher than Li, and Li is higher than Be.

cross-section formation in spallation reactions resulting from the bombardment of protons on atoms of C and O (Figs. 2 and 3). The link is best shown through phase-space arguments. In the high-energy region, the breakup of the excited nuclei into a given configuration is proportional to the number of possible channels, which is itself a function of the binding energy, and also of the number of bound excited states for this configuration. Above one hundred MeV or so, the cross sections reach a plateau, which they maintain all the way up to the highest energies. As expected, the cross section for the formation of ^{9}Be is the smallest, while Li and B isotopes have rather similar values.

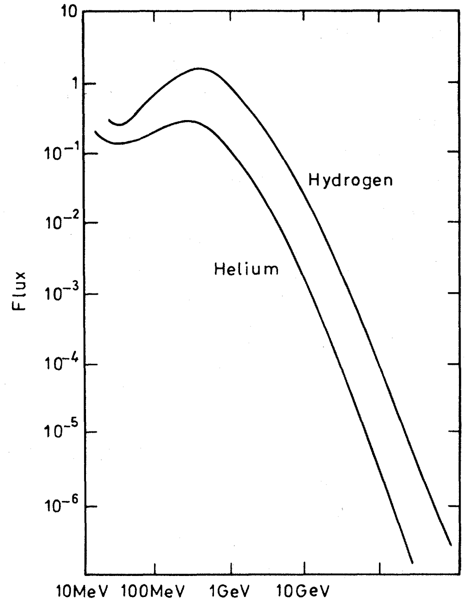

FIG. 4. Galactic cosmic-ray fluxes ($cm^{-2}s^{-1}ster^{-1}$ MeV^{-1}) of H and He as a function of energy per nucleon. Most of the particles are found between 100 MeV and 2 GeV.

A similar sequence is observed in the cosmic-ray fluxes of the Li, Be, and B isotopes (Figs. 4 and 5). It is indeed believed that these fast nuclei are generated when high-energy C and O collide with interstellar protons. This process, together with its counterpart (fast protons hitting interstellar C and O), is believed to be the main source of ^{6}Li, ^{9}Be, ^{10}B, and ^{11}B. The arguments for this view will be summarized later. Let us note for the moment that this process meets our requirement: the light nuclei are not subjected to high temperatures after their formation.

III. OBSERVATIONS AND EXTRAPOLATIONS

While most people agree on the observed abundances of the light elements (see Wilson and Rood, 1994), there are often appreciable differences of opinion regarding the uncertainties to be attached to these numbers. In this article these uncertainties will be dealt with at two levels: first at a "reasonable" or "optimistic" level, and second at an "extreme" or "pessimistic" level. This second level is preferred when major issues are at stake, as, for instance, the possible existence of dark baryonic matter or nonbaryonic matter.

A. Observations of deuterium

This discussion is based on the analysis of Geiss (1993). An argument based on the measurements of the helium isotopic ratio in the solar system can be used to give an estimate of the D/H ratio in the galactic gas at the birth of the sun, 4.5 billion years ago. In the meteorites, noble gases include two components, called "solar" and "planetary." The first component is interpreted as captured solar wind. It should represent, with minor modifications, the present composition of the sun. The second component is believed to have been already present in the protosolar gas. It provides information on the isotopic composition of the interstellar cloud from which the solar system was formed.

The helium isotopic ratio in the present solar surface [from solar wind measurements, from the solar component of gas-rich meteorites, and from the moon (Geiss *et al.*, 1970; Jeffrey and Anders, 1970; Black, 1972)] is $(^3\mathrm{He}/^4\mathrm{He}) = 4.50 \pm 0.4 \times 10^{-4}$. This ^{3}He is understood as being the sum of the presolar (planetary) component, $^3\mathrm{He}/^4\mathrm{He} = 1.5 \pm 0.3 \times 10^{-4}$ (Eberhardt, 1974), plus a component resulting from the fusion of D into ^{3}He in the early days of the formation of the sun. [The D/H ratio in the solar wind is indeed less than 3×10^{-6} (Epstein and Taylor, 1972). A subtraction of the two components should then yield an estimate of the D/H ratio in the protosolar nebula (Black, 1971; Reeves, 1971; Geiss and Reeves, 1972).]

The evaluation of the ^{3}He/H is made through the estimated solar value of ^{4}He/H of 0.10 (± 0.01), from observations of solar flares (Heasley and Milkey, 1978), or from solar model fitting (Bahcall and Ulrich, 1988; Turck-Chièze *et al.*, 1988), corresponding to a helium mass fraction of 0.28 (± 0.02). One obtains $^3\mathrm{He/H} = (\mathrm{D} + ^3\mathrm{He})/\mathrm{H}$ (protosolar) $= 4.1(\pm 1.0) \times 10^{-5}$. The corresponding D/H obtained by subtraction of the planetary component is

$$\mathrm{D/H\ protosolar} = 2.6 \pm 1.0 \times 10^{-5} .$$

Deuterium is also observed in molecular form on many planets, where it has been subjected to a number of physico-chemical alterations (Geiss and Reeves, 1981). On Jupiter and Saturn, however, these processes were probably of minor importance. The original solar system abundance ratio is best obtained from their atmospheres. Encrenaz and Combes (1982) estimate the range of D/H in Jupiter to be between 1.2 and 3.1×10^{-5}, and Gautier and Owen (1989) between 2 and 4×10^{-5}. These data are quite consistent with the analysis from the helium isotopic ratio.

The Hubble satellite has obtained the D/H ratio in the present interstellar gas in front of Capella (Linsky *et al.*, 1992). The ratio is $\mathrm{D/H} = 1.65(+0.07, -0.18)10^{-5}$. This value agrees well with values obtained along other lines of sight (Vidal-Madjar, 1991; McCullough, 1992), but the quoted uncertainties are much reduced. From the previous discussion it appears that the D/H depletion between the birth of the sun and now is probably less than a factor of two.

B. Observations of ^{3}He

The protosolar value is obtained from the planetary component of helium:

$$^3\mathrm{He}/^4\mathrm{He(protosolar)} = 1.5 \pm 0.3 \times 10^{-4} .$$

Observations of ^{3}He in the present interstellar medium have been reported by Rood *et al.* (1984), through the detection of the 3.46 cm hyperfine line of $^3\mathrm{He}^+$. The measurements range from $^3\mathrm{He/H} = 4 \times 10^{-5}$, in the giant HII region W43, to 8×10^{-5} in W51, and to 15×10^{-5} in W3. An upper limit of 2×10^{-5} has been obtained in W49 and M17S. In view of the fact that ^{3}He is also generated in stars—it has been observed at a value of $^3\mathrm{He/H} = 10^{-3}$ in a planetary nebula (Bania *et al.*, 1993; Rood *et al.*, 1992)—this dispersion may not be too surprising. The upper limit of 2×10^{-5} in W49 and M17S is quite within the range of the solar system estimates.

C. Primordial D and (D + ^{3}He) values

We have no data on D and ^{3}He before the birth of the solar system. Meaningful comparisons with the BBN yields imply that we can correctly extrapolate the observed abundances all the way back to the moment of primordial nucleosynthesis. This extrapolation involves poorly known properties of stellar and galactic physics and, in this sense, remains uncertain and model dependent.

We can take advantage of the fact that, in view of its low destruction temperature, deuterium is only destroyed by stars. As a simple rule, the fractional surviving amount of D at any one time is given by the fractional amount of the original galactic gas that has never been incorporated into a star (astrated) at that time.

The case of ^{3}He is more complicated, since this nucleus is both produced (by the $\mathrm{D} + p \Longrightarrow {}^3\mathrm{He}$ reaction) and destroyed (by the $^3\mathrm{He} + {}^3\mathrm{He}$, or $^3\mathrm{He} + {}^4\mathrm{He}$ reactions) in stars. The best strategy (Yang *et al.*, 1984) is to study independently the D/H ratio and the $(^3\mathrm{He} + \mathrm{D})/\mathrm{H}$ ratio. Since the burning of D generates ^{3}He, the sum of both is less likely than D itself to be model dependent. On the other hand, the fact that, in contrast to D, ^{3}He is a normal product of stellar nucleosynthesis is an added complication to the extrapolation of the second ratio.

Several groups (Delbourgo-Salvador, 1985; Brown, 1992; Steigman and Tosi, 1992; Vangioni-Flam, Olive, and Prantzos, 1994) have tried to determine the primordial values of D and (D + ^{3}He) from galactic evolution models based on our present knowledge of factors such as the star formation rate, the initial mass function of stars, the evolutionary abundance curves of elements such as O, C, N, Fe, etc. There is general agreement that, at solar birth, the remaining D/H was more than 25 percent of its primordial value and the $(\mathrm{D} + {}^3\mathrm{He})/\mathrm{H}$ more than 40 percent. Consequently the upper limits on the primordial D/H and $(\mathrm{D} + {}^3\mathrm{He})/\mathrm{H}$ can both be taken as $< 10^{-4}$. For the lower limits I adopt for $\mathrm{D/H} > 1.6 \times 10^{-5}$ and $(\mathrm{D} + {}^3\mathrm{He})/\mathrm{H} > 3 \times 10^{-5}$.

How much confidence should we put in these extrapolations? Their main weakness is that they are based on a rather naive conception of galaxies in their youth, one on which recent observations cast serious doubts. When the first stellar generation was formed (redshift z larger than 4 or 5), the universe was at least one hundred times denser than today and galactic collisions are likely to have played an important role, deeply influencing the rate of star formation and of nucleosynthesis. The assumption that galaxies are isolated systems quietly transforming their gas into stars is difficult to reconcile with the picture of galactic merging and consequent starbursts recently observed (Djorgovski, 1987; Carlberg, 1990; Guiderdoni and Rocca-Volmerange, 1990a).

Most galactic evolution models are implicitly based on the Eggen *et al.* (1962) model of the disk formed by the collapse of the primitive halo. Recently, however, models incorporating the occurrence of galactic collision (mergers) have been discussed, and attempts have been made to evaluate the effects of these processes on the galactic evolution of element abundances (Guiderdoni and Rocca-Volmerange, 1990b; Colin and Schramm, 1992). Clearly the extrapolation to the big bang brings us to a chapter of astrophysics about which we know very little: the birth and early evolution of galaxies.

To take this added uncertainty into account, we shall arbitrarily divide the evolution of the galaxy into two periods: an early period corresponding to metallicity less than one-tenth solar and a recent period with metallicity larger than one-tenth solar.

1. Early galactic era

The near constancy (within a factor of two) of the ^{7}Li/H ratio, (see Fig. 7 below) in old stars with heavy-element abundances (the astronomer uses the term "metallicity") ranging from one-thousandth to one-tenth the solar value can be used as an argument *against* important depletion of fragile elements in the galactic gas during this period. [These stars, formed very early in the life of the galaxy, are called Population II (Pop II) stars. Later-formed stars, with higher metallicity, are called Pop I stars. Metallicity is usually based on the solar value.] Although D burns at a lower temperature than ^{7}Li, as far as astration is concerned, this difference is not significant. This argument limits to a factor of two the depletion of D/H in the very early days of the galaxy.

(One could argue that the "constancy" of ^{7}Li/H in this period may be due to a balance between destruction by astration and formation by cosmic rays. However, the rate of growth of Be and B in the same period does not appear to support this objection.)

2. Recent galactic era

Here the stellar and nuclear processes are better understood. The estimates of D/H$<10^{-4}$ and (D+^{3}He)/H$<10^{-4}$ from the "conventional wisdom" galactic models discussed earlier will be adopted as reasonable estimates for this later phase. Compounding the effects of both periods, we can set the extreme limits of (D+^{3}He)/H and D/H both at $<2\times10^{-4}$.

D. Observations of ^{4}He

^{4}He is generated in the big bang and also by stars. The observed abundances range from 23 percent to 30 percent in fractional mass, showing the gradual effect of stellar energy generation in galaxies. The search for the primordial yield is conducted in objects that have been least contaminated by stellar synthesis, such as the so-called extra galactic HII regions (blue compact galaxies), as is evidenced by their comparatively low abundance of heavier isotopes. One strategy consists of a plot of He versus O or C or N, which is tentatively extrapolated to zero value for these stellar-generated isotopes. Another estimate is obtained by averaging on the objects with the lowest metal abundances. (Fig. 6; Kunth and Sargent,

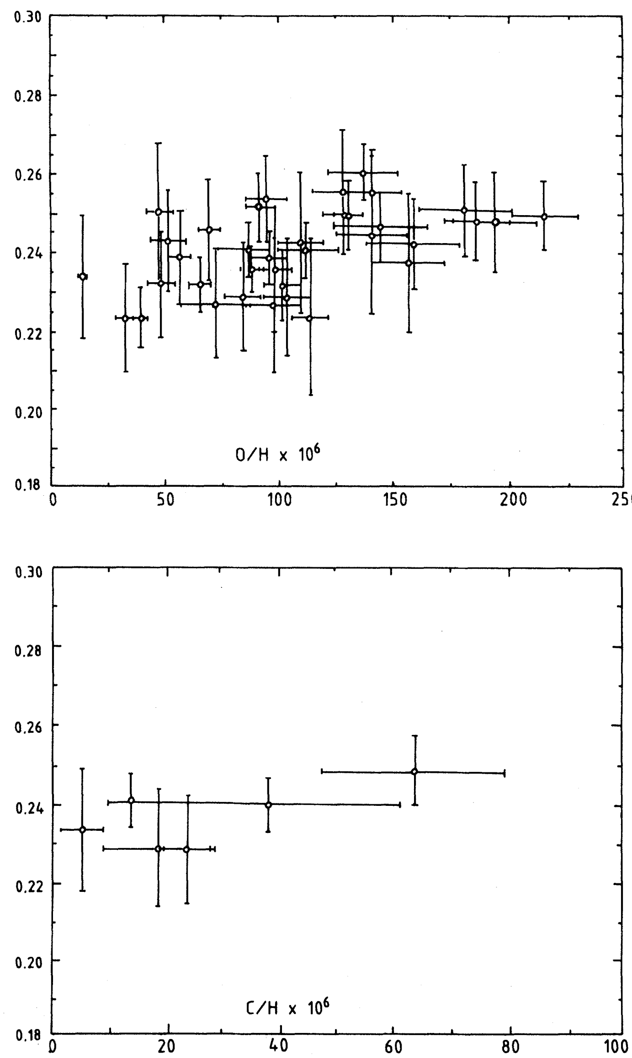

FIG. 6. Mass fraction Y of ^{4}He as a function of the abundance ratios O/H, C/H, and N/H in galaxies with very small metal abundance.

1983; Pagel, 1989; Campbell, 1992; Fuller *et al.*, 1991; Pagel *et al.*, 1992).

The data support a value of Y (the helium mass fraction)$\approx$0.23. There is much debate regarding the uncertainties attached to this number. To quote a few, we find $0.226 < Y < 0.248$ (Kunth and Sargent, 1983); $0.224 < Y < 0.254$ (Boesgaard and Steigman, 1985); $0.215 < Y < 0.245$ (Steigman, Gallagher, and Schramm, 1989); $0.22 < Y < 0.24$ (Walker *et al.*, 1991); $0.21 < Y < 0.24$ (Smith, Kawano, and Malaney, 1993); $0.223 < Y < 0.232$ (Pagel, 1993); $0.228 < Y < 0.238$ (Campbell, 1992).

It has often been suggested (Schramm, 1993) that, following the example of the nuclear physicists, astronomers should distinguish between statistical uncertainties (where here are evaluated as ± 0.01) and systematic uncertainties (more difficult to pin down but perhaps as large as ± 0.01). In view of the important conclusions regarding dark matter to be drawn from the helium-4 abundance, I shall adopt as "reasonable" limits $0.22 < Y < 0.24$. Taking into account the recent low value of Terlevich *et al.* (1992) and following the advice of

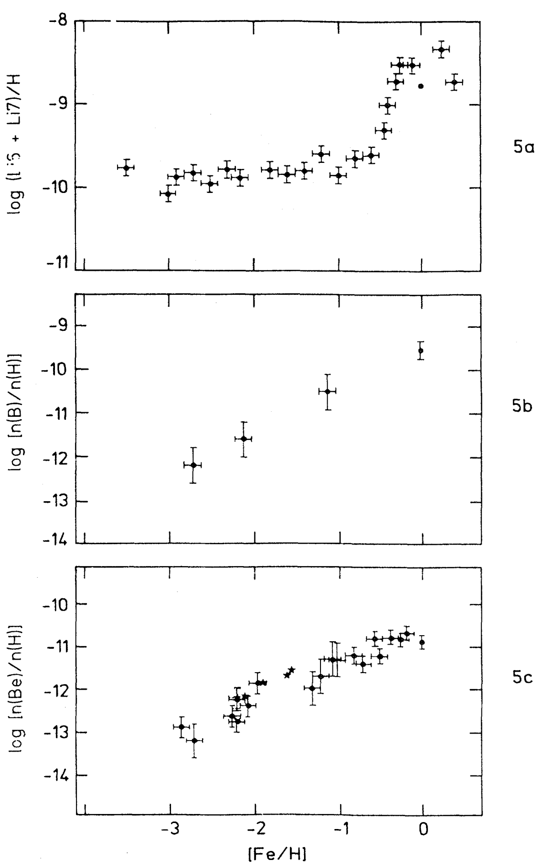

FIG. 7. Stellar abundances of the elements Li, B, and Be in the old (Pop II) stars as a function of their iron content (normalized to solar; "-3" means "10^{-3} of the solar iron abundance"). Note the Spite plateau in the Li/H curve, the signature of the big-bang contribution. No such plateau has yet been found for Be and B. If found, these plateaus would represent evidence for a first-order quark-hadron phase transition.

Pagel (private communication), I shall set the extreme limits at $0.21 < Y < 0.25$.

Helium emissivity calculations by Smits (1991) have recently been shown by Skillman *et al.* (1993a) to increase the previously reported helium abundances by as much as 4% (or 0.010). The same authors (Skillman *et al.*, 1993b) report a helium mass fraction of 0.239 ± 0.006. These results give weight to the view that systematic errors should be considered and that uncertainties are usually underestimated.

E. Observations of ^{7}lithium

The best argument in favor of a big-bang contribution to the abundance of ^{7}Li comes from the fact that, for stars with very small abundances of the heavy elements, the lithium abundance, at the level of $^7\mathrm{Li/H} \approx 1.2 \times 10^{-10}$, is independent of the metallicity (the Spite plateau in Fig. 7; Spite and Spite, 1982a, 1982b; Spite *et al.*, 1992; see also Hobbs and Duncan, 1987; Rebolo *et al.*, 1988). Nuclei generated during thte galactic lifetime have abundances that increase gradually with metallicity.

Like ^{4}He and ^{3}He, but unlike D, ^{7}Li is not a pure big-bang child. It is produced by stars, as well, and by the galactic cosmic-ray (GCR) bombardment of interstellar gas. These contributions are believed to be responsible for the gradual increase of lithium abundances to a Li/H ratio of $(1-2) \times 10^{-9}$ (Magazzu, Rebolo, and Pavlenko, 1992; Martin, Rebolo, and Pavlenko, 1992; Spite and Spite, 1993) as the stellar metallicity increases from one-tenth solar to the solar value (Fig. 7).

The value of the lithium isotopic ratio is known for the protosolar gas, 4.5 billion years ago ($^7\mathrm{Li}/^6\mathrm{Li} = 12.5 \pm 0.5$; Krakowski and Muller, 1967) and also for the present galactic gas ($^7\mathrm{Li}/^6\mathrm{Li} = 12.5 \pm 4.0$; Lemoine *et al.*, 1993a, 1993b). However, Meyer *et al.*, 1992, report a smaller value). Smith, Lambert, and Nissen (1992) have reported the detection of ^{6}Li at the level of $^6\mathrm{Li}/^7\mathrm{Li} = 0.05 \pm 0.02$ in a metal-poor (one-hundredth of solar) Pop II star.

To obtain the big-bang yield from these observations we still have to estimate the lithium depletion from surface processes in these old stars. This turns out to be a difficult problem. Tentative indications may be obtained by applying a simple rule: the depletion should be at least equal to, but not much more than, the dispersion (by a factor of two) of the measured abundances.

Several groups (Vaulclair, 1988a, 1988b; Delyannis *et al.*, 1990; Charbonnel, 1992; Pinsonneault *et al.*, 1992) have considered this problem. Effects such as atomic and turbulent diffusion processes, meridional circulation, rotational mixing, and gravity waves have been explored. After an extensive discussion of these mechanisms, Michaud and Charbonneau (1991) conclude that "the observed Li abundance has probably been reduced by a factor of two, so that the original abundance was probably equal to $^7\mathrm{Li/H} = 3 \times 10^{-10}$." They add, however, that "more calculations are needed to better establish this value." Following this discussion, I shall adopt the values $1.0\times10^{-10}<{}^7\mathrm{Li/H}<3\times10^{-10}$ as reasonable estimates.

Three observations can be used to set a strong upper limit to the primordial abundance of ^{7}Li/H. First, the measurement of the lithium isotopic ratio in one Pop II star (Smith, Lambert, and Nissen, 1993) has been used as an argument against important lithium depletion by rotational mixing as advocated by Pinsonneault *et al.* (1992). An assumed depletion of ^{7}Li by a factor of ten would imply a ^{6}Li depletion irreconcilable with the observation of a ^{6}Li/^{7}Li of five percent. Second, the measurement of ^{7}Li/^{6}Li in the interstellar medium by Lemoine *et al.* (1993a, 1993b) can be used to assess the contribution of stars (probably asymptotic giant branch stars, as discussed by Scalo, 1986) to the present abundance of ^{7}Li/H. The computations of Reeves (1993) leave little room for the possibility of a BBN contribution larger than $^7\mathrm{Li/H}\approx10^{-9}$. Third, supernova 1987A has provided an upper limit for the abundance of Li in the Large Magellanic Cloud. This limit has been used to argue against a primordial ^{7}Li/H as large as 10^{-9} (Baade and Magain, 1988; Sahu *et al.*, 1988), although the uncertainties are large (Malaney and Alcock, 1989).

For all these reasons, one can adopt an initial lithium value of $^7\mathrm{Li/H}<10^{-9}$ as an extreme upper limit.

The mean lithium abundance in unevolved Pop I stars is $\mathrm{Li/H}=2\times10^{-9}$, with an uncertainty of a factor of two. This number agrees well with the meteoritic data. Lithium has been observed in a large number of stars in various states of evolution. The study of the abundance variations is a very, important tool for the elucidation of the physics of stars (Charbonnel, 1992; Michaud and Charbonneau, 1992) and galactic evolution (Audouze and Tinsley, 1976; Audouze *et al.*, 1983; Abia and Canal, 1988; Fuller *et al.*, 1988; Vangioni-Flam *et al.*, 1990; D'Antona and Mateucci, 1991; Brown, 1992).

F. Observations of beryllium

Beryllium has been observed in 38 main-sequence Pop I stars by Boesgaard (1976). The mean value is $\mathrm{Be/H}=1.31\pm0.36\times10^{-11}$, in good agreement with the solar value $\mathrm{Be/H}=1.4\times10^{-11}$ measured by Chmielewski *et al.*, 1975. The meteoritic (carbonaceous chondrite) value is slightly higher, $\mathrm{Be/H}=2\times10^{-11}$, but well within the uncertainties (Grevesse and Noels, 1993).

Like lithium, the element beryllium shows a systematic decrease with decreasing stellar metallicity for stars with less than one-tenth of the solar metallicity (Fig. 7). Unlike Li, however, Be shows no sign of a plateau at very small metallicity. At a metallicity of one-thousandth of the solar value, it has decreased by a factor of one hundred from its Pop I value ($\mathrm{Be/H}\approx10^{-13}$); (Ryan *et al.*, 1990; Gilmore *et al.*, 1991, 1992; Gilmore, 1993; Rebolo *et al.*, 1993; Boesgaard and King, 1993), while lithium reaches its plateau ($\mathrm{Li/H}\approx10^{-10}$) at one-tenth of its Pop I ($\mathrm{Li/H}\approx10^{-9}$) value.

G. Observations of boron

The ratio B/H has been measured in 16 Pop I stars (*A* and *B* type) by Boesgaard and Heacox (1978). The mean value is $\mathrm{B/H}=2\times10^{-10}$ with an uncertainty of "a factor of two or three." For the sun, Kohl *et al.* (1977) measure $\mathrm{B/H}=4\ \times10^{-10}$, again with a factor-of-two uncertainty. For the solar system (carbonaceous chondrites) Anders and Grevesse (1989) give the value $\mathrm{B/H}=7.5\pm0.6\times10^{-10}$, almost four times larger than the mean stellar value quoted above. The reason for this discrepancy is not clear. Anders (private communication) offered the following explanation: "boron is volatile in steam (e.g., volcanic furnaroles), and it is known that Cl chondrites were exposed to H_2O vapor." This exposure could have brought extra boron into the chondrites, thereby accounting for their apparent overabundance of this element. Until this question is elucidated, it appears wiser to rely on the stellar average quoted above and to select $\mathrm{B/H}=2\times10^{-10}$ for Pop I stars.

[It is worth mentioning that *if* the boron meteoritic value is the correct present galactic abundance, we would need an extra source of boron (10 and 11). The GCR contribution, as evaluated by the Be abundance and the ratio of the spallation cross sections, would be two or three times too small.]

In Pop II stars, boron behaves just as beryllium does, but the number of stellar measurements is much smaller (Duncan *et al.*, 1992). With decreasing metallicity, it decreases to $\mathrm{B/H}=10^{-12}$ with no sign of a plateau.

The boron isotopic ratio in the solar system is $^{11}\mathrm{B}/^{10}\mathrm{B}=4.05\pm0.05$ (Shima and Honda, 1962). We have no other measurement of this ratio.

IV. THE CASE FOR BIG-BANG NUCLEOSYNTHESIS

The best argument in favor of big-bang nucleosynthesis is the fact that the abundance ratio to hydrogen of four nuclei—D, ^{3}He, ^{4}He, and ^{7}Li—is reasonably well accounted for by adjusting only one parameter, the baryonic density ρ_b. The required value of ρ_b falls within the "window" determined by astronomical observations.

Baryonic densities ρ_b are usually given in terms of the critical density ρ_c between an open and a closed universe: $\Omega_b\equiv\rho_b/\rho_c$. The critical density has the value $\rho_c=3H_0^2/8\pi G_N=1.89\times10^{-29}\ (H_0/100)\ \mathrm{g\,cm^{-3}}$, where H_0 is the Hubble constant in km/sec/megaparsec. The fact that the value of H_0 is known only within a factor of two (between 50 and 100 in these units) introduces one major source of uncertainty in the whole of cosmology. As a useful number, the "most probable" value $H_0=75$ km/sec/megaparsec ($2.2\times10^{-18}\ \mathrm{sec^{-1}}$ in more standard units) gives $\rho_c\approx10^{-29}\ \mathrm{g/cm^3}$.

The lower limit of ρ_b is given by the density of luminous matter (stars and hot gas). It corresponds to $\Omega_b\approx0.005$ (Tremaine and Biney, 1987). A strong upper limit to the total density of all matter ($\Omega<3$) is suggested by the lack of observed deceleration of galaxies (Sandage, 1987). Dynamical studies of the rotation curves of galaxies or stability arguments from clusters of galaxies give estimates of $\Omega \approx 0.1$ to 0.2. As we shall see presently, the best value of Ω_b from big-bang nucleosynthesis is well within the window of astronomical requirements.

It is worth mentioning at this point that, before the first CERN results in 1981, the BBN model had two *free* parameters, the second one being the number N of families of elementary particles (Shvartsman, 1969; Steigman, Schramm, and Gunn, 1977; Yang *et al.*, 1984). Based on the observed abundance of helium, the BBN model predicted that this number N should be three or at most four (more about this later). The latest results of LEP (Adeva *et al.*, 1992; $N=3.05\pm0.05$) provide a brilliant confirmation of the BBN prediction.

Another argument in favor of big-bang nucleosynthesis comes from the fact that no celestial objects are known to have a helium fraction of less than 22% to 23%. This universal lower limit strongly suggests a process acting on a cosmic scale. It is in fact a good measure of the helium abundance in highly metal-poor galaxies, which are consequently believed to have undergone very little stellar nucleosynthesis. No other processes are known which could generate He without also generating oxygen and iron.

The Spite plateau for the lithium abundance in very old stars constitutes still another argument in favor of big-bang nucleosynthesis. The fact that the lithium abundance levels off to a value where it is independent of stellar metallicity is an argument against a stellar nucleosynthesis or GCR spallation origin of this component of lithium and consequently an argument in favor of BBN.

The observation of deuterium at the level of $D/H \approx 10^{-5}$ is also an argument for big-bang nucleosynthesis. As mentioned in the nuclear physics overview, in view of its extreme fragility the stellar synthesis of D/H comes out many orders of magnitude below this value. The second best candidate, GCR production by proton spallation of ^4He, comes out at a level of $D/H=10^{-8}$ at best.

One can look at big-bang nucleosynthesis in the following way. From the Hubble observations of galaxies, we know that the universe is expanding. From the Einstein general relativity equations applied to a homogeneous universe, we know that is also cooling, and we can relate the rate of cooling to the present matter density. This tells us that the universe was warmer and denser but does not tell us *how high a temperature* it reached in the past.

The discovery of the fossil 3 K radiation can be interpreted as a proof that it reached at least 3000 K (in fact, much more) when this radiation was last scattered. In the same vein, if we assume that the universe has been at a temperature as high as 10^{10} K (weak-interaction decoupling), then we find a simple explanation for the observed abundance of D, ^3He, the metal-poor galactic abundances of ^4He, and the Pop II abundance of ^7Li. In the calculation, the required baryonic density is in the range allowed by astronomical observations, and the required number of families of elementary particles is consistent with accelerator experimentation. No other cosmological theory has come anywhere near doing as well.

A. Physical parameters of the big-bang theory

Of major importance for an understanding of primordial nucleosynthesis is the temperature behavior of the various time scales involved in the physics of the expansion. We distinguish the macroscopic time scale related to the expansion from the various microscopic time scales related to the particle reaction rates.

In the early universe, the Einstein equations yield the simple relation

$$(\dot{R}/R)^2 = 8\pi G_N \rho/3 \ , \tag{1}$$

where R is the scale factor, $\dot{R}$ its time derivative, and ρ the total energy density; G_N is Newton's constant. Dimensionally we have

$$t(\text{expansion}) \propto (G_N\rho)^{-1/2} \ . \tag{2}$$

The energy density ρ is dominated by the set of all relativistic particles "i" ($kT \gg M_i c^2$) for which $\rho_i \propto g_i T_i^4$, where g_i is the statistical multiplicity factor ($g_i = 2$ for the photons). A "demographic factor" g^* is usually introduced to represent the combined effect of all the particle species, bosons (b), and fermions (f), on the mass density of the early universe:

$$\rho = (g^*/2)\rho(\text{photons}) \propto g^* T_\gamma^4(\text{photons}) \tag{3}$$

where T_γ is the temperature of the photon gas. For reasons to be discussed shortly, the various relativistic gases may not be at the same temperature. This fact will influence their density contributions and will appear in the expression for g^* in the following way:

$$g^* = \Sigma g_{bi}(T_{bi}/T_\gamma)^4 + \tfrac{7}{8}\Sigma g_{fi}(T_{fi}/T_\gamma)^4 \ . \tag{4}$$

The factor $\frac{7}{8}$ reflects the difference between the statistics for bosons (b) and fermions (f). The multiplicities are given by g_{bi} and g_{fi}. At one MeV, for instance, in the Standard Model of elementary particles, the "standard" demography consists of photons, electrons, positrons, three types of neutrinos, and their antineutrinos, all at the same temperature ($T_{bi} = T_\gamma = T_{fi}$), thus $g^*(T=1\ \text{MeV}) = 2 + \frac{7}{8}\times 10 = 10.75$. Note that because it is related to the number of species with $M_i c^2 < kT$, g^* is a function of temperature.

The cosmic expansion time scale can be written (see Fig. 8) as

$$t(\exp) \propto (g^* G_N)^{-1/2} T^{-2} \ . \tag{5}$$

For the reaction time scale, let us consider a reaction of the type $A + B \Longrightarrow C + D$, for instance, the capture of a neutrino by a neutron to give an electron and a proton. The capture cross sections σ are a function of the energy with a given power (usually positive). The probability of one capture event is given by the product of the cross section times the relative velocity, $\langle \sigma v \rangle$, averaged over the velocity distribution of particles at temperature T. This average value is proportional to the strength of the interaction times a power m of the temperature. For the weak interaction implied in the neutrino capture reactions, $m=2$ and the Fermi constant G_F appears squared:

$$\langle \sigma v \rangle(\text{weak}) \propto G_F^2 T^2 \ . \tag{6}$$

The probability of reaction per unit volume is proportional to $\langle \sigma v \rangle$ times the number density of capturing particles per unit volume $n(T)$, which, in the expanding universe, is $n(T) \propto T^3$. Thus the lifetime t(reac) for a given neutrino (in our example) to interact with neutrons (the inverse of the reaction probability) is given by

$$t(\text{reac})(n+\nu \Longrightarrow p+e) \propto G_F^{-2} T^{-5} \ . \tag{7}$$

For other reactions, the temperature exponent will be $(-m-3)$, which, for all physical processes of importance, is larger than the exponent value of 2 characteristic of the expansion time scale. The result is that, in the very early universe, the reaction time scales are always shorter than the expansion time scale. The two curves meet (Fig. 8) at the *decoupling temperature* T_d, which is a function of many parameters such as the ratio of the coupling constants (G), the demography of the universe (through g^*), and other factors influencing the cross section, including its energy dependence (through the power m).

One important consequence of this comparative behavior of the time scales is that all the physical processes are in *statistical equilibrium* at early moments of the expansion. The relative abundances of the reacting particles are then given by laws of mass action, such as Boltzmann's law, and are independent of past situations. After decoupling, in contrast, the processes occur in a state of disequilibrium and the abundances reflect past history.

For the weak interactions the decoupling temperature is obtained by equating the two time scales [Eqs. (2) and (5)]:

$$T_d \propto (g^*)^{1/6}(G_N)^{1/6}(G_F)^{-2/3} \approx 1 \text{ MeV} \ . \tag{8}$$

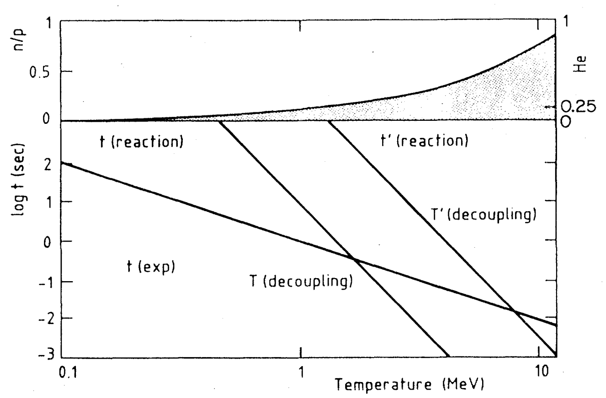

FIG. 8. Decoupling of the neutrino interactions. The abscissa gives the cosmic temperature, and the ordinate gives the age of the universe in seconds. The curve t(exp) gives the relation between cosmic age and temperature $[t(\text{exp}) \propto 1/[(G_N)^{1/2}T^2]]$ in the standard big-bang model. The t(reaction) curve is the mean reaction time for weak interactions involving neutrino capture and emission, $t(\text{reaction}) \propto 1/(G_F^2 T^5)$. Cosmic time runs from right to left. At temperatures below the crossing of the curves $(T < T_D \approx 1$ MeV), the neutrino interactions are too slow to keep pace with the expansion, and the neutron-proton equilibrium abundance is no longer assured. A particle interacting with a G value smaller than the Fermi constant $(G < G_F)$ would have its reaction-time curve t'(reaction) shifted to the right in the diagram, leading to a higher decoupling temperature T'_D. On the upper part of the digram, the n/p ratio is shown as a function of temperature. The Boltzmann-equilibrium equation is shown as the border of the shaded area. It ceases to apply at the decoupling temperature T_D. On the scale at the right, one can read at T_D an approximate value of the BBN helium yield. The position of the T_D can be altered by changing G_F or G_N, as may be seen from the expressions for the time scales. The BBN model turns out to be a very sensitive test of the "constancy" of the coupling constants.

B. The first seconds

Many crucial events take place during the first seconds of the universe: (a) the weak-interaction decoupling event at one MeV at one second or so (actually the μ and the τ neutrinos decoupled at a slightly higher temperature than the e neutrinos); (b) the electron-positron annihilation around 0.5 MeV at four seconds or so; (c) the nucleosynthesis of the light nuclei around 0.1 MeV at one hundred seconds or so. A useful approximation for the time-temperature relationship is $t(\text{sec}) = 1/T^2(\text{MeV})$.

At $T > 1$ MeV the weak-interaction equilibrium is ensured by the reactions

$$n+\nu_e \Longleftrightarrow p+e; \ \nu_e+p \Longleftrightarrow n+e; \ n \Longleftrightarrow p+e+\nu_e \ .$$

In consequence, the neutron-proton ratio is given by Boltzmann's law (top of Fig. 8):

$$n/p = \exp(-\Delta M/kT) \ , \tag{9}$$

where ΔM is the neutron-proton mass difference (1.293 MeV). As T decreases, the n/p ratio goes down from one to a value close to 0.2 at decoupling. After T_d the neutrino capture reactions are too slow to maintain the weak-interaction equilibrium.

The electron-positron annihilation at 0.5 MeV creates a flux of new photons, which increases slightly the photon radiation (more exactly, they lower the cooling rate). Since the neutrino interactions are now very weak, the neutrinos are essentially decoupled from the rest of the universe, and they receive no share of the energies released by the annihilation. As a result, at this moment, the neutrino temperature T_ν becomes slightly lower than the photon T_γ. This effect can be calculated through the conservation of entropy per covolume during the annihilation phase. The entropy density is proportional to the number density of relativistic interacting particles. To compute this effect we define $g^*_{\rm int}$ as the "interaction demographic factor": it contains only the species that are in physical interaction at a certain temperature (here $T \approx 0.5$ MeV). Here i and f will stand for the initial state (before annihilation) and the final state (after annihilation):

$$\begin{aligned} &s_i = g_i^*{}_{\rm int} T_i^3 = s_f = g_f^*{}_{\rm int} T_f^3 : \\ &g_i^*{}_{\rm int} = [2 + \tfrac{7}{8}(2+2)] \quad (\text{photons} + \text{electrons}) , \\ &g_f^*{}_{\rm int} = [2](\text{photons only}) , \\ &\text{thus } T_f / T_i = (11/4)^{1/3} = 1.40 . \end{aligned} \tag{10}$$

Since the neutrinos did not receive their share of this annihilation phase, they remained at T_i while the rest of the universe reached the temperature T_f. (Fields *et al.*, 1992, have discussed small corrections to this effect.) Today we measure a photon temperature T_γ of 2.7 K. Thus we expect that the neutrino radiation is at $T_\nu = 1.9$ K. Such a cosmological background of neutrinos in a necessary consequence of the theory if indeed the universe has reached temperatures over one MeV in the past. Because of the low mean energy of these particles (one meV), the detection of a neutrino background is outside the realm of contemporary technology.

The previous example explains why various relativistic gases can be at different temperatures as expressed in Eq. (4) for g^*. Suppose, for instance, that there exist right-handed neutrinos, interacting as a weak-interaction particle, but with a coupling constant $G' \ll G_F$. If their decoupling temperature, evaluated through Eq. (8), is larger than 107 MeV (the mass of the muon), they will receive no share from muon annihilation at the equivalent temperature. Conservation of entropy during this phase, evaluated as in the previous example, will allow a determination of the temperature of this radiation, still lower today than the neutrino temperature. This example will be helpful later on.

Around 0.1 MeV the gamma rays (the tail of the Bose-Einstein photon energy distribution) are no longer numerous enough to keep the deuteron population in statistical equilibrium with the nucleons. This is the onset of primordial nucleosynthesis. Through the reaction $n + p \Longrightarrow \gamma + \mathrm{D}$, the population of D increases rapidly, as shown in Fig. 9. As they reach a ratio of D/H of 10^{-3} or so, they undergo further nuclear reactions and are transformed into mass-3 nuclei: $\mathrm{D} + p \Longrightarrow {}^3\mathrm{He} + \gamma$; $\mathrm{D} + n \Longrightarrow T + \gamma$. The population of these mass-3 nuclei increases in turn, as the D decreases.

The same fate befalls the mass-3 nuclei as ^{4}He starts its rise through the ${}^3\mathrm{He} + {}^3\mathrm{He} \Longrightarrow {}^4\mathrm{He} + 2p$. Because there are no stable nuclei with masses 5 and 8, the ^{4}He suffers essentially no further nuclear depletion. Only a very small fraction of its population gets transformed into ^{7}Li through the reactions ${}^4\mathrm{He} + T \Longrightarrow {}^7\mathrm{Li} + \gamma$ and ${}^4\mathrm{He} + {}^3\mathrm{He} \Longrightarrow {}^7\mathrm{Be} + \gamma$, followed, after many days, by

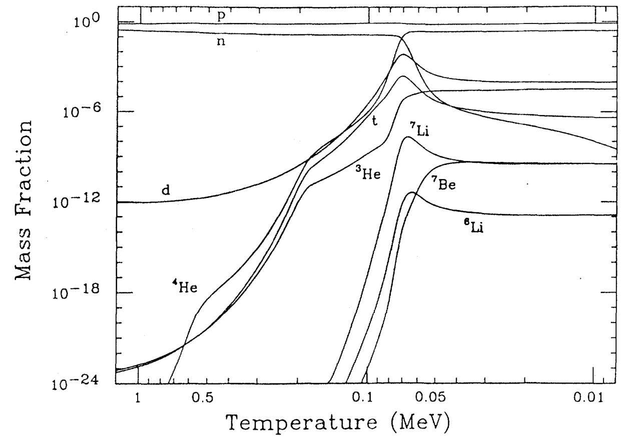

FIG. 9. Time history of primordial nucleosynthesis. The mass fractions of the various nuclei are displayed as a function of temperature in billion degrees at the top, or as a function of time in seconds at the bottom (from Smith, Kawano, and Malaney, 1993).

${}^7\mathrm{Be} + e \Longrightarrow {}^7\mathrm{Li} + \nu_e$. Some ^{7}Li is further destroyed but, because of its larger Coulomb charge ($Z=4$), the ^{7}Be remains almost intact (Fig. 9).

The relative abundances of the light nuclei generated during this period of cosmic nuclear activity are related to two key parameters of the physical conditions during expansion: (a) the neutron-to-proton ratio and (b) the nucleon-to-photon ratio during the active nuclear phase (T between 1 and 0.01 MeV).

The n/p ratio is related to the decoupling T_d and hence to the value of g^* [Eqs. (4), (8), and (9)]. In the standard BBN model, the value of g^* is fixed by the assumed demography of the universe. We shall discuss later the implications of models with different values of g^*.

Let us focus our attention on the fate of the neutrons after decoupling. They may either beta-decay (with a lifetime of 890±4 seconds, Mampe *et al.*, 1989) or interact with a proton to form a D. The probability of this last result is proportional to the density of nucleons (baryons) ρ_b. At low ρ_b the neutrons beta-decay; at higher ρ_b they undergo nuclear reactions and are essentially all processed to ^{4}He (with very minor formation of the light nuclei with mass 2, mass 3, and mass 7). For baryonic densities smaller than the critical density, the yields of other nuclei are negligible.

In order to identify the density-temperature profile of cosmic matter during primordial nucleosynthesis, we assume that no important entropy-generating processes have taken place from $T=0.1$ MeV until now. (This point will be discussed again later.) Thus the nucleon-to-photon ratio should have remained constant. This number is usually characterized by the baryonic number, $\eta = n$ (baryons)/n (photons). The strategy would be to find the value of η today and to use it in order to fix the nucleonic density-temperature profile in the past. The number of photons is obtained from the fossil radiation

[n(photons)$\approx$400 cm^{-3} for T=2.736]. Since we have no independent knowledge of the baryonic density, η is left as a parameter.

For all densities greater than the density of luminous matter ($\eta > 3\times10^{-11}$) most of the neutrons present at decoupling find their way into a ^{4}He nucleus. Thus the abundance of helium is related to the n/p ratio at the weak charged-interaction decoupling temperature by $Y\approx 2(n/p)/(1+n/p)$. It is a fair monitor of the value of the weak decoupling T_d [through Eq. (9)] and hence of the value of g^* [through Eq. (8)]. On the other hand, it is only weakly dependent on the baryonic number η, as shown in Fig. 10(c).

The abundance of D, on the other hand, depends strongly upon the baryonic number. At high ρ_b, the fractional abundance of D surviving destruction by p or n capture to produce mass-3 nuclei becomes very small. For instance, if the baryons were at the critical density, the D/H ratio would be 10^{-12}, seven orders of magnitude below the observed values. The mass-3 nuclei show similar behavior but somewhat less pronounced. The behavior of mass-7 nuclei is more complex, with a hill (from ^{7}Li formation), a valley, and a raising slope (from ^{7}Be) at higher η [Fig. 10(d)].

The yields are computed with the help of a reaction network incorporating a large number of reaction cross sections for the production and destruction of the various nuclei. They were first computed in a seminal paper by Wagoner *et al.* (1967). The nuclear data were regularly reviewed by Caughlan and Fowler (1988). The subject has been thoroughly reanalyzed recently by Krauss and Romanelli (1990); Thieleman *et al.* (1991); Walker *et al.* (1991); and Smith, Kawano, and Malaney (1993); see also Kawano (1992).

C. Comparison with observations: (a) homogeneous big-bang nucleosynthesis

The comparison of theoretical yields with observations will be carried out in two steps. First, in the context of

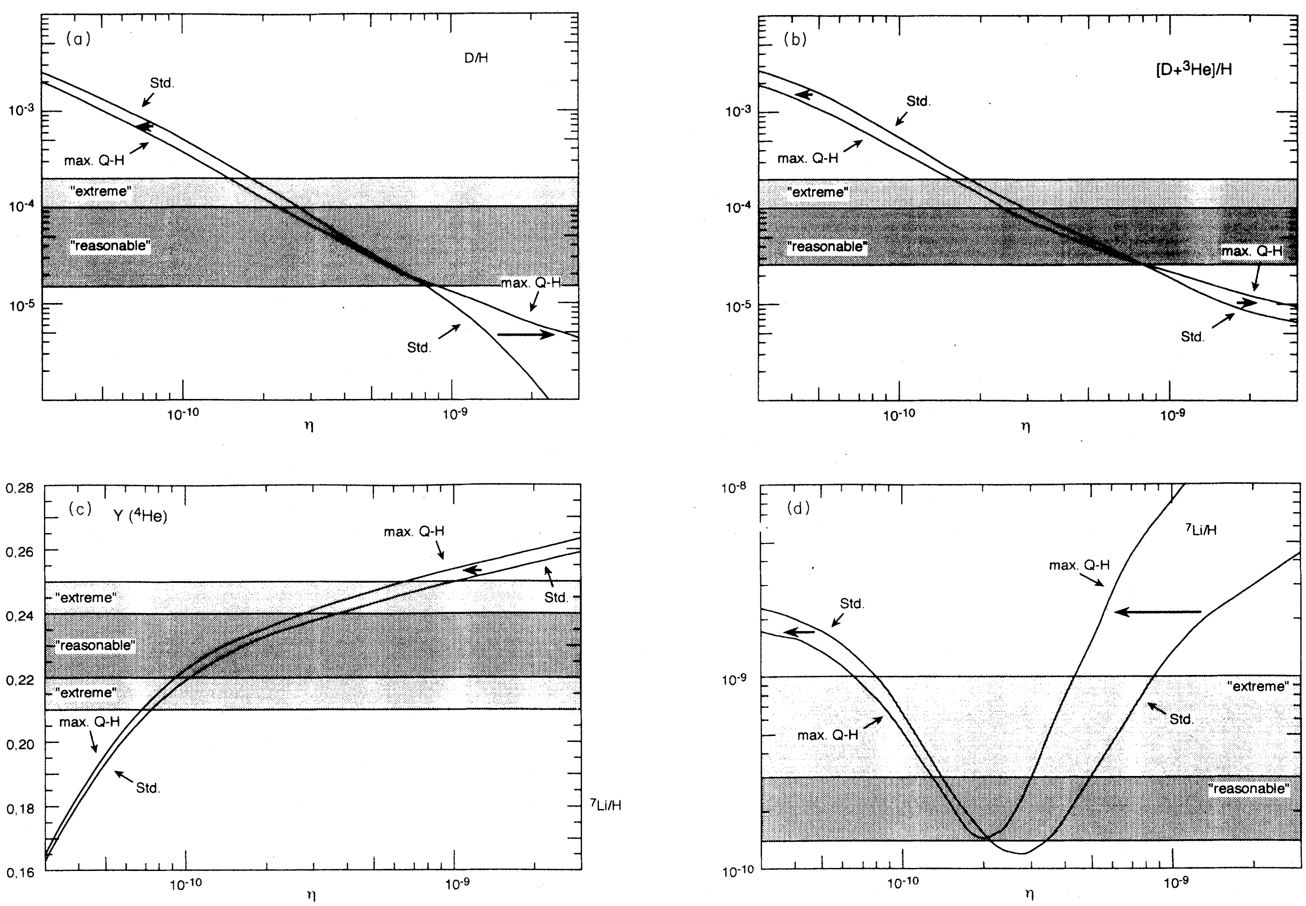

FIG. 10. Light nuclei yields in the early universe. The abscissa is logη=n(nucleons)/n(photons). The curves labeled STD refer to homogeneous density BBN. The curves labeled "Max Q-H" are obtained in the case $d < 10^{5.0}$ cm. For ^{7}Li the curve Max Q-H applies at all values of d. The arrow indicates the direction in which the quark-hadron effects displace the yield curves when the assumed contrast parameter R is varied from 1 to 10^4. Dark-shaded areas are obtained from the reasonable limits. The full areas (light+dark) correspond to the extreme limits.

the so-called "standard BBN model," we shall assume that, at $T=1$ MeV, cosmic energy density was homogeneously distributed in space. In the second section, the potential effects of the quark-hadron phase distribution at $T\approx 150$ MeV will be discussed. By creating inhomogeneities, both in energy density and in the proton-to-neutron ratio, this phase could have generated yields differing from the homogeneous state.

Corresponding to a given value of η_{10} (the baryonic number in units of 10^{10}), the baryonic density, $\rho_b = 0.6\eta_{10}\times 10^{-31}$ g/cm^3, expressed in units of the closure density ($\Omega_c = 1$), is given by

$$\Omega_b(Ho/100)^2 = 3.7\times 10^{-3}\eta_{10} \tag{11}$$

where Ho is the Hubble constant. With the value $Ho=75$ km/s/megaparsec, we have $\Omega_b = 6.6\times 10^{-3}\eta_{10}$. The range of possible Ho introduces an uncertainty of a factor of two on each side.

The lower limit of $\mathrm{D/H} > 1.6\times 10^{-5}$ gives an upper limit of $\eta_{10} < 8$. The "reasonable" upper limit of $< 10^{-4}$ corresponds to $\eta_{10} > 3.0$, and the "extreme" upper limit of (D) $< 2\times 10^{-4}$ corresponds to $\eta_{10} > 2.0$. The results on (D+3He)/H are practically the same. For ^{4}He we have selected "reasonable" limits at $0.22 < Y < 0.24$ corresponding to $1.2 < \eta_{10} < 4.0$ and "extreme" limits of $0.21 < Y < 0.25$ corresponding to $0.8 < \eta_{10} < 12$. For ^{7}Li our "reasonable" value $1.0\times 10^{-10} < {}^7\mathrm{Li/H} < 3.0\times 10^{-10}$ gives $1.4 < \eta_{10} < 5.0$, while the "extreme" limit of ${}^7\mathrm{Li/H} = 10^{-9}$ corresponds to $0.80 < \eta_{10} < 9.0$.

Thus, in the standard BBN model, for the *upper limits,* ^{4}He and ^{7}Li gives a similar reasonable value of $\eta_{10} < \approx 5$, while all four nuclei give a similar extreme upper value of $\eta_{10} < \approx 10$. For the *lower limits,* the strongest reasonable value comes from D at $\eta_{10} > 3$, while Y and ^{7}Li only set $\eta_{10} > 1$. In the same fashion, the strongest extreme value comes from D at $\eta_{10} > 2.0$, while Y and ^{7}Li only set $\eta_{10} > 0.8$. The situation is illustrated in Fig. 10, parts (a)–(d).

Putting together these requirements of the various nuclei, we have for the standard case the reasonable limits $2 < \rho_b(10^{31}$ g/cm$^3) < 3$, while for the extreme limits $1.2 < \rho_b(10^{31}$ g/cm$^3) < 6$.

In terms of the critical density, we find $0.01 < \Omega_b < 0.08$ and the extreme limits of $0.008 < \Omega_b < 0.15$. We realize that the uncertainties on Ho are now the main source of uncertainties on Ω_b.

D. Comparison with observations: (b) inhomogeneous big-bang nucleosynthesis—the quark-hadron phase transition

Around 1980, progress in particle physics called into question the "standard model" used up to then to compute BBN yields. Witten (1984) showed that the quark-hadron (Q-H) phase transition, occurring at $T\approx 150$ MeV, could have generated density inhomogeneities, leading after weak-interaction decoupling at $T=1$ MeV, to inhomogeneities in the neutron-to-proton ratio during big-bang nucleosynthesis at $T=0.1$ MeV.

One key point is the order of Q-H transition (or transitions, since there are both a chiral and a deconfinement transition, taking place at approximately the same energy). If the transition is of second order, no inhomogeneities are created, and the standard model is acceptable. If it is of first order, the degree of inhomogeneity depends upon a number of parameters, including the duration of the following overcooling period. The effects on BBN yields may or may not be important.

Information on the parameters of the phase transition can, in principle, be obtained both from high-energy collisions of heavy nuclei in the laboratory and from QCD calculations on a lattice. On the laboratory front, it appears likely that the quark-gluon soup has been obtained at CERN, but so far very little information has been reliably extracted from these experiments (Morel, 1992). Events showing the presence of the particle J/Ψ, expected if the transition has really occurred, have indeed been observed, but other explanations for these events cannot be excluded. As yet, there is no clear signal that the quark-gluon soup has really been observed. The hope is that the lead-lead collisions, programmed at CERN in a year or two, will give pertinent experimental results on these processes.

Despite the fact that QCD calculations on a lattice have been pursued for a number of years, we are still far from definite results. They still involve too many simplifying assumptions to warrant convincing conclusions. One interesting new development is the fact that, while previous QCD computations took into account only the presence of gluons, in the last two years quarks have been progressively introduced. From the (apparently nonlinear) volume dependence of one of the parameters of the transition. Fukugita and Hogan (1991) have concluded that the transition is not of first order. However, according to many specialists, this analysis is far from being conclusive, and the order of the transition must still be considered as uncertain (Gottlieb, 1991; Petersson, 1993). The phase transition is most likely a weakly first-order transition (Martinelli, 1994).

A second important effect of the inclusion of quarks in the computation is a lowering of the estimated temperature of the transition. While in the pure gauge computations (quark free) the temperature was found to be around 220–250 MeV. It falls well below 200 MeV when quarks are taken into account (Toussaint, 1992). In the case of one heavy and two light quarks, it is close to 150 MeV. The lower the temperature, the higher the expected density contrast in the baryon inhomogeneities potentially created by the transition (if indeed it is of first order). This is of importance, since the yield differences between homogeneous and inhomogeneous big-bang nucleosynthesis increase with the contrast parameter.

The surface tension of the nucleated bubbles of hadrons in the sea of quarks is another important parameter of the transition. It is usually quoted in terms of the ratio s/T_c^3. If s/T_c^3 is small, nucleation will start just below the critical temperature, leading to the formation of many small bubbles. The mean distance between the bubbles will also be small. Inversely, if s/T_c^3 is large the nucleation process will be delayed, leading to the formation of fewer bubbles at large interbubble distances. Brower *et al.* (1992) have evaluated this quantity with QCD computation on a lattice including quarks. They obtain an upper limit of $s/T_c^3 < 0.12$.

It appears advisable to treat the uncertainties on the effect of the Q-H phase transition as uncertainties on the BBN yields. While in the case of the homogeneous big bang there is only one free parameter, the ratio η of the number of nucleons to the number of photons, the Q-H phase transition introduces three new parameters.

The first is the effective density contrast R between the high-density regions (the clumps) and the low-density regions. The word "effective" takes into account the fact that the contrast created at the beginning of the nucleation process [proportional to $\exp(-M(\text{proton})/kT_c)$] is later amplified by the percolation process (the bubbles growing to occupy the whole of space) to several tens of times its initial value (Alcock *et al.*, 1987). The second is the fraction of the mass f_v in the clumps (the clumpiness). The computations of Fuller *et al.* (1988) suggest that $f_v < 0.1$.

The third is the average distance d between the clumps. A convenient unit is the present value of d in light hours (lh). One lh today corresponds to $10^{4.4}$ cm at $T = 1$ MeV and approximately one meter at the Q-H phase transition, when the horizon scale was approximately ten km. At large values of d ($d > 10^3$ lh) the neutron could not diffuse before nucleosynthesis. Computations made on the assumption of large values of d would give the same results as computations based on a density-inhomogeneous standard model (i.e., with no n/p inhomogeneities). At the lower end of the scale, $d < 0.01$ lh, proton diffusion becomes important, the clumpiness is partly erased, and we fall back on the homogeneous model.

As discussed before, the value of d is mostly a function of the surface tension of the nucleated bubbles. An approximate expression has been given by Kurki-Suonio (1991):

$$d \approx 1.0 \text{ lh } (T_c/200 \text{ MeV})^{-2}(s/T_c^3)^{3/2} \tag{12}$$

(a numerical mistake, reproduced in Reeves, 1990, has been corrected).

With the most recent estimate of $T_c = 140$ MeV and the upper limit given for the surface tension, d should be less than 0.1 lh. In this range the effect of the Q-H transition on the BBN yields are small but not negligible.

After the early computations of Applegate and Hogan (1985), Malaney and Fowler (1988), and Reeves *et al.* (1988), calculations, including neutron diffusion during big-bang nucleosynthesis, have been made covering the whole parameter space corresponding to the uncertainties on the value of these parameters (Alcock, Fuller, and Matthews, 1987; Applegate *et al.*, 1988; Turner, 1988; Terasawa and Sato, 1989; Kurki-Suonio *et al.*, 1990; Reeves, Richer, *et al.*, 1990; Terasawa, 1993; Malaney and Mathews, 1993). N. Terasawa has kindly made his most recent results available for this review. His latest improvements include large mesh size (100), large nuclear networks (up to $14N$), stretching function method for spatial zoning, and updated nuclear reactions rates, especially for neutron-rich unstable nuclei.

Many interesting features have come out of these recent systematic computations (Reeves and Terasawa, 1994). For instance, one finds that the upper value of η_{10} corresponding to the upper limit of the ^{7}Li/H ratio is systematically *decreased* by the effects of the Q-H transition (i.e., at all values of d and f_v). Thus the upper limits of η_{10} obtained from the homogeneous case ($\eta_{10} < 5$ in the reasonable limit; $\eta_{10} < 9$ in the extreme limit) are robust limits. They can be said to be "Q-H proof" [see Fig. 10(d)].

Next we use the results of Brower *et al.* (1992) on the value of the bubble energy density to set a reasonable limit to the interbubble distance ($d < 10^{3.8}$ cm today). In this d range, the Q-H effects do not alter the upper limits on η_{10}. Thus we retain both the reasonable upper limit ($\eta_{10} < 5$) and the extreme upper limit ($\eta_{10} < 10$). Unfortunately the lower limits to η_{10} are never "Q-H proof." We have to consider each case in detail. From D/H and (D+3He)/H we have a reasonable lower at $2.0 < \eta_{10}$ and an extreme lower limit at $1.5 < \eta_{10}$. For ^{4}He we have a reasonable lower limit at $0.8 < \eta_{10}$ and an extreme limit at $0.5 < \eta_{10}$. For ^{7}Li/H we have a reasonable lower limit at $1.0 < \eta_{10}$ and an extreme limit at $0.5 < \eta_{10}$.

In a nutshell, including the uncertainties on the Q-H phase transition, the "reasonable" limits of D/H and (D+^{3}He)/H correspond to $2.0 < \eta_{10} < 9$ and the "extreme" limits to $1.5 < \eta_{10} < 9$. For ^{4}He the "reasonable" limits correspond to $0.8 < \eta_{10} < 4.0$ and the "extreme" limit to $0.5 < \eta_{10} < 12$. For ^{7}Li the "reasonable" limits correspond to $1.0 < \eta_{10} < 5.0$ and the "extreme" limit to $0.5 < \eta_{10} < 9.0$. These new limits are shown in Figs. 10(a)–(d).

Putting together these requirements of the various nuclei, we have the reasonable limits $2.0 < \eta_{10} < 5.0$, or $1.2 < \rho_b(10^{31}\ \text{g/cm}^3) < 3$, while for the extreme limits $1.5 < \eta_{10} < 10$ or $0.9 < \rho_b(10^{31}\ \text{g/cm}^3) < 6$. In terms of the critical density, we find for the reasonable limits $0.008 < \Omega_b < 0.08$ and for the extreme limits $0.005 < \Omega_b < 0.15$. It is worth recalling that the extreme upper limit is fixed by all four nuclei. For ^{7}Li it is independent of uncertainties on the parameters of the Q-H transition. For the other nuclei it depends only on the assumption that $d < 0.1$ lh, which seems well established (Brower *et al.*, 1992). The lower limits are based only on the values of D and ^{3}He, extrapolated from the protosolar values to the time of big-bang nucleosynthesis. They are functions of the parameters of the Q-H transition. In this sense they are not as safe as the upper limits.

The following conclusions emerge from these calculations.

(1) The limit $\Omega_b < 0.08$ is reasonably certain, while $\Omega_b < 0.15$ is a very strong upper limit. The universe is definitely not closed by baryonic matter; $\Omega_b < 1$.

(2) At the density level required by dynamic arguments ($\Omega \approx 0.1$ to 0.2), no strong case can be made for the existence of nonbaryonic matter.

(3) There is probably, but not certainly, a fair amount of dark baryonic matter ($\Omega_b > \Omega_{\text{luminous}} \approx 0.005$).

(4) The possibility of BBN production of Be and B in a strongly inhomogeneous universe has been explored by Malaney and Fowler (1988), Boyd and Kajino (1989), and Thomas *et al.* (1993). Besides the fact that the computations reported earlier do not support the hypothesis of strong homogeneities, there are several objects to this view. On the observational side, no B or Be plateau at low stellar metallicity has yet been established, which would substantiate the idea of a BBN contribution to these elements. On the theoretical side, the computed B/Be ratio ≈ 100 (Terasawa, 1992) appears to be in conflict with the observed ratio at low metallicity (≈ 10).

We are in position to answer the question: what are the effects of the Q-H phase transition on big-bang nucleosynthesis? As seen from Fig. 10, they are small at best. They do not change the upper limits to the baryonic numbers but may decrease the lower limits (reasonable and extreme) by $\sim 40\%$.

V. BIG-BANG NUCLEOSYNTHESIS AS A TEST BENCH

The success of the BBN model in accounting for the abundances of the four cosmological nuclei has served as a test for new ideas in particle physics. In this view, new theories must not disturb the reasonable agreement obtained with the simple version, which is compatible with already experimentally tested particle physics.

A. New weakly interacting particles?

The crucial parameter is the demographic factor g^* defined in Eq. (4). The yields of big-bang nucleosynthesis (in particular helium-4) are controlled by the value of g^* at weak-interaction decoupling g^* ($t_{\text{ex}} = t_{\text{reac}}$). In the standard BBN model the temperature is $T_{\text{D}} \approx 1$ MeV.

The good agreement between the evaluation of the number of families of elementary particles obtained through the abundance of helium and through experiments at LEP has already been mentioned. It is important to note that the two methods do not test exactly the same things. The LEP method puts a limit on the number of channels leading to the decay of the Z_0 particle, the boson carrier of the weak interaction. This limit leads to a limit on the number of left-handed neutrinos (Fermi interaction) with less than half the mass of the Z_0 ($M < 45$ GeV). It is hardly sensitive to hypothetical particles with interactions much weaker than the Fermi interaction.

Conversely, the value of g^* in big-bang nucleosynthesis includes, in some way, all existing particles (Steigman, 1992). The respective roles of hypothetical species are a function of their mass (i.e., whether they are relativistic or not at decoupling) and also of the strength of their interactions, through the fact that their relative number density is affected by this strength.

Bertolini and Steigman (1992) discuss the relevance of the success of the BBN model for limiting the number of possible weakly interacting fermions as a function of their coupling constant ($G < G_F$). Such weakly interacting particles would decouple at a temperature T'_{D} higher than ordinary fermions (Fig. 8). Entropy conservation requirements [Eq. (10)] specify that the density ratio $\propto (T_{fj}/T\gamma)^4$ of a fermion j at big-bang nucleosynthesis be given by

$$(T_{fj}/T\gamma)^4 = \{g^*_{(\text{int})}\ (\text{at } T = 1\ \text{MeV})/g^*_{(\text{int})}\ (\text{at } T_{\text{D}fj})\}^{4/3}\ . \tag{13}$$

As an example, Bertolini and Steigman (1992) argue that if the helium mass fraction Y were shown to be less than 0.24, the decoupling temperature of one assumed extra fermion would have to be more than 150 MeV, which would require its coupling to be less than 10^{-5} of the Fermi coupling. Its exchange boson would have a mass larger than 10^{12} eV.

B. The leptonic number of the universe

Our universe is strongly asymmetric with respect to baryons. The number of baryons appears to be much larger than the number of antibaryons, $n_b \gg n_b^*$, and the number of baryons is much smaller than the number of photons, $n_b/n_\gamma \approx 3 \times 10^{-10}$. What is the situation in the leptonic world? We have good reasons to believe that the number of neutrinos is closely equal to the number of antineutrinos. And both numbers are close to the number of photons.

Big-bang nucleosynthesis can give some information on these matters (Reeves, 1972). The lepton number, $L_l = \Sigma_i (n_i - n_i^*)/n_\gamma$ (where Σ_i is the sum over all the leptons) is conserved during cosmic expansion. To explore the situation, BBN yields are computed with L_l as a free parameter, and the compatibility with the observed abundances is tested. These computations show that the effect of L_l on BBN is negligible unless $L_l > 10^{-2}$, so that the present contribution of electrons ($Le \approx 0.85 L_b \approx 10^{-10}$, for electric neutrality) can be left out. We consider only neutrinos.

Lepton number is best discussed in terms of the ratio of the chemical potential μ_ν to the thermal energy of the neutrinos, $\xi_\nu \equiv \mu_\nu / kT_\nu$, a number which is conserved during the expansion. The neutrino energy distribution is given by

$$dN_\nu \propto E_\nu^2 dE_\nu / \{(\exp[(E_\nu/kT_\nu)-\xi_\nu])+1\} \ . \quad (14)$$

The antineutrino ν^* energy distribution is given by the same expression, with $\xi_{\nu*}=-\xi_\nu$. In terms of ξ_j, the leptonic number of a species of neutrino j is given by $L_j \approx 0.075 \{F_2(\xi_j)-F_2(-\xi_j)\}$, where

$$F_2(\xi_j)=\int_0^\infty x^2 dx[\exp(x-\xi_j)+1]^{-1} \quad (15)$$

or

$$L_j \approx 0.25\{(\xi_j)+0.1(\xi_j)^3\} \ . \quad (16)$$

The effect of ξ_j on the energy density of the degenerate neutrino species j is given by

$$\rho(\nu_j)+\rho(\nu_j^*)=aT_j^4/c^2\{\tfrac{7}{8}+\tfrac{15}{4}(\pi^{-2})\xi_j^2+\tfrac{15}{8}(\pi^{-4})\xi_j^4\} \ , \quad (17)$$

where "a" is Boltzmann's constant (Beaudet and Goret, 1976). For our purpose it will be useful to express this density in terms of the ratio of neutrino to photon energy density,

$$\rho(\nu_j)+\rho(\nu_j^*)/\rho_\gamma = \tfrac{7}{8}(T_j/T_\gamma)^4\{1+\tfrac{30}{7}(\xi_j/\pi)^2+\tfrac{15}{7}(\xi_j/\pi)^4\} \ . \quad (18)$$

A nonzero value of L_l (corresponding to degeneracy in the neutrino species) would influence big-bang nucleosynthesis in two different ways.

(a) From Eqs. (16) and (17) it implies a higher number density than a zero value of L_l and hence it increases the expansion rate over the standard gT^4 contribution of each species (this applies to the three types of neutrinos).

(b) In the case of the electron-neutrinos, a further role is played by their influence on the neutron/proton ratio. Degeneracy of the electron-neutrinos implies a population difference between the ν and ν^* which, in turn, shifts the balance of the $n+e^+ \Longleftrightarrow p+\nu_{e*}$ and $p+e^- \Longleftrightarrow n+\nu_e$ reactions in such a way that the new equilibrium ratio is given by

$$n/p=\exp\{\xi_e+\Delta M(n-p)/kT\} \ . \quad (19)$$

For example, if $\xi_e>0$, then $L_e>0$ [Eq. (16)] and $n(\nu_e)>n(\nu_{e*})$. Compared to the nondegenerate case ($\xi_e=0$), the equilibrium will be shifted toward more protons and fewer neutrons, as seen in Eq. (19). Remember that the n/p ratio at decoupling governs the ^{4}He abundance.

BBN yields have been calculated (David and Reeves, 1980a, 1980b) with three parameters: the usual η (the nucleonic density); ξ_e (the neutrino-electron degeneracy), and an effective parameter ξ_μ playing the role of ξ_j in Eq. (18) and representing the combined density effects of degeneracy in the three types of neutrinos. The computations show that the effect on the yields of an increase of ξ_μ can be compensated for by an increase in ξ_e. Compatibility with the observations of cosmological nuclei can be achieved for the range of matter density going from the luminous matter lower bound all the way up to a critical density of baryons. *In other words, the BBN model alone cannot exclude a universe closed by baryons if appropriate neutrino degeneracies are assumed.*

The situation is different, however, if another cosmological constraint is taken into account, in relation to the problem of galaxy formation (Steigman, 1986). In the standard scenario of galaxy formation (gravitational growth from early density fluctuations), the time delay required to obtain a full-blown galaxy corresponds to a redshift difference of at least 10^3 (and more likely 10^4) between the moment when the universe became matter-dominated and today. Indeed, galaxies cannot start condensing before the end of the radiative era, i.e., before the moment when the radiation (photons plus neutrinos) density became equal to the matter density (Meszaros, 1974). Since the radiation density is proportional to T^4, while the matter density is proportional to T^3, the ratio of the temperature at equal densities of radiation and matter ($T_{\rm eq}$) to the present temperature of the fossil radiation ($T=2.74$ K) is given by the ratio of the densities today:

$$T_{\rm eq}/2.74=(1+z_{\rm eq})=\rho_{\rm matter}/\{\rho(\nu_j)+\rho(\nu_j^*)+\rho_\gamma\} \ , \quad (20)$$

where$=z_{\rm eq}$ is the redshift at equality. The condition $z_{\rm eq}>10^3$ corresponds to $\rho(\nu_j)+\rho(\nu_j^*)/\rho_\gamma<10$ if $\Omega_b=1$. This, in turn, would require [from Eq. (18)] $\xi_\mu<5$. But at $\Omega_b=1$, BBN chemical abundance compatibility would require $\xi_\mu>20$! Thus $\Omega_b=1$ is excluded.

Combining the $z_{\rm eq}>10^3$ for galaxy formation with the BBN requirements, we get an upper limit of $\xi_\mu<3$, corresponding to $\Omega_b<0.07$ (with $Ho=75$). The calculations also determine the maximum range of $-0.15<\xi_e<0.3$; limiting the degeneracy of the electron neutrinos and implying no strong matter/antimatter asymmetry in the leptonic world. The corresponding range of L_e is $-0.04<L_e<0.08$.

From the physics point of view, most grand unified theories (but not all) have comparable values of L_b and L_l, so that we should expect $L_l \ll 1$. This appears to be in agreement with the prediction of the BBN model. (Note: In this discussion all neutrinos were assumed to be massless. The masses are unimportant if $m_\nu<0.3$ meV. See Steigman, 1986.)

One word of caution, however. In view of the difficulty of building a coherent theory of galaxy formation solely on the assumption of linear gravitational growth from primordial density fluctuations, as exemplified by the latest COBE results, the possibility of other scenarios, including accelerating condensation factors or late-time phase transitions (Schramm, 1993) cannot be ruled out at this point. Thus the constraint $z_{\rm eq}>10^3$ should perhaps not be considered as definitely established.

The effects of possible neutrino oscillations between the different flavors have been considered by Savage *et al.* (1991) and Entqvist *et al.* (1990). They may further increase the uncertainties on the leptonic numbers.

C. The constancy of the "constants" of physics

The success of the BBN model in accounting in a satisfactory way for the observed and extrapolated abundances of the quartet of cosmological nuclei can also be used to test constancy of the laws of physics which are instrumental in fixing these abundances. The theory of the big bang is naturally based on the assumption that the laws of physics and their "constants" are really constants. The idea here is to investigate how much variation can be tolerated in the numerical values of these constants, while still retaining satisfactory agreement between the computed yields and the observations.

In a sense, this argumentation may appear circular: we assume the constancy of the laws to test the validity of the big bang as a scientific theory, and then we use the theory to test the constancy of the constants. However, the fact that there are other proofs for these constancies may be taken as complementary evidence in favor of the argumentation (Bahcall and Schmidt, 1967; Wolfe, Brown, and Roberts, 1976; Irvine, 1983; Reasenberg, 1988).

As mentioned earlier, the primordial abundance of helium is closely related to the n/p ratio at the weak-interaction charge decoupling temperature T_D given by Eqs. (8) and (9). Variation in any of the terms of the expression of T_D would result in a different helium abundance (Yang *et al.*, 1984). To remain within our extreme limits, the Newton constant G_N should have varied by less than 25% and the Fermi constant G_F by less than 6% (assuming of course no correlated variations).

An interesting study of the meaning of possible variations in the Fermi coupling constant has been presented recently by Scherrer and Spergel (1992). G_F is, in fact, the ratio of two physical parameters: the true gauge coupling constant g_f of the electroweak SU(2) interaction and M_W, the mass of the W boson carrying the interaction: $G_F=(2)^{1/2}g^2/8M_W^2$. But since, in the gauge formalism, the mass of the M_W is itself the product of g_f^2 times the expectation value of the Higgs field $\langle\Phi\rangle^2$, the value of G_F is essentially independent of g_f and can be related directly to the expectation value of the Higgs field: $G_F=1/(2)^{1/2}\langle\Phi\rangle^2$. The "present" value of $\langle\Phi\rangle$ is 246 GeV. The present limits on helium do not allow a variation of more than three percent of its quantity in the last 15 billion years.

Barrow (1987) has discussed various ways in which modifications of the fundamental constants could affect the helium yield. In addition to the effect of modified G_N and G_F in Eq. (9) one should also consider the fact that the neutron-proton mass difference is probably due in large part to the electromagnetic interaction, thus being sensitive to any modification of the fine-structure constant. Besides its undoubted influence, as well, on this mass difference, a modification of the strong-coupling constant would also alter the very weak binding energy of the deuteron and hence influence the production of the isotopes D and ^{3}He (Pochet *et al.*, 1991).

Cosmologies with extra geometric dimensions (i.e., superstring theories) have received much attention in recent years. A popular version involves ten dimensions, thus adding six new compact dimensions, over and above the familiar three space and one time dimensions. The radius of curvature of these extra dimensions would be of the order of the Planck length (10^{-33} cm), far smaller than the smallest dimensions within reach of presently operating accelerators (the accelerators at CERN or at Fermilab can probe to $\approx 10^{-18}$ cm). Energies of the order of the Planck mass (10^{19} GeV) would be required to excite the corresponding modes. In these cosmological models, the values of the coupling constants of the various forces depend upon the radius of curvature of these compact dimensions. If, as is the case in our familiar 3D expanding universe, these radii are changing with time, the coupling constants would also vary. The success of the BBN model and the implied constancies of the constants of physics also imply a high stability for the curvatures. This stability can be used to place severe constraints on the choice of acceptable superstring theories (Kolb, Perry, and Walker, 1986).

D. Limits on global entropy increase since big-bang nucleosynthesis

The entropy density of the fossil radiation is inversely proportional to the nucleon-to-photon ratio η. A lower limit to η in the present universe is given by the number density of "shining" matter divided by the number density of photons in the cosmological radiation: $\eta_{\rm now} > 5\times 10^{-11}$. An upper limit to this number at primordial nucleosynthesis is obtained from our extreme limit $\eta_{\rm BBN} < 10^{-9}$ from the data discussed earlier.

The ratio of these numbers, $\eta_{\rm BBN}/\eta_{\rm now} < 20$, gives the maximum increase due to all possible reheating of the universe between BBN and the present (Vauclair *et al.*, 1993). This ratio is proportional to the third power of the ratio of the temperatures after and before reheating: $T_{\rm after}/T_{\rm before} > 2.7$. Cosmologically significant entropy increases are often associated with phase transitions leading to episodes of inflation. The maximum consequent increase in the scale factor $R_{\rm after}/R_{\rm before} = \exp(Ht) < 2.7$, hence $Ht < 1$, where $H=(8\pi G\rho_v/3)^{1/2}$ where ρ_v is the corresponding vacuum energy. This value of Ht gives the limiting possible effect of inflationary episodes after bigband nucleosynthesis.

One possible suggestion for such a phase transition may come from solar physics (Hill *et al.*, 1990). The analysis of the solar data suggests that the MSW (Mikheyev, Smirnov, and Wu) effect (Mikheyev *et al.*, 1986) may be at work if two of the neutrinos have a (square) mass difference of some 10^{-4} to 10^{-7} eV^2, corresponding to individual neutrino masses of a fraction of an electron volt. It has been suggested that such masses could be involved in a phase transition related to family symmetry of elementary particles.

VI. THE CASE FOR GALACTIC COSMIC-RAY ORIGIN

In contrast to ^{7}Li, ^{9}Be is not an important secondary product of hydrogen burning, either in stars or during primordial nucleosynthesis. It is a typical case of a nucleus for which we know of no low-energy production mechanism.

The rate of formation of ^{9}Be in GCR reactions is given by the product of the flux of high-energy protons (≈ 16 cm^{-2}sec^{1}) times the cross sections for ^{9}Be formation by proton collision on the most abundant targets, ^{16}O and ^{12}C, (≈ 5 mb) times the abundance ratio of these targets to hydrogen ($\approx 10^{-3}$) in space. The approximate equality between the product of the formation rate times the age of the galaxy, on the one hand, and the beryllium-to-hydrogen ratio in recent stars ($\approx 10^{-11}$) is the best evidence for a major GCR contribution to some of the light elements.

Comparing the ratios of the spallation cross sections of protons on O and C to the ratios of the stellar abundances of Li, Be, and B confirms the view that some of the light nuclei are generated by galactic cosmic rays. The analysis, to be detailed later, shows that the GCR mechanism can satisfactorily account for the nuclei ^{6}Li, ^{9}Be, and ^{10}B; it gives a major contribution to the abundance of ^{11}B and a minor (10%) contribution to the abundance of ^{7}Li (Reeves *et al.*, 1970; Meneguzzi *et al.*, 1971; Reeves *et al.*, 1973; Austin, 1981; Walker *et al.*, 1985; Arnould and Forestini, 1989).

A. Physical parameters of the galactic cosmic-ray nucleosynthesis

In the nuclear sector, the important parameters are the excitation functions for spallation reactions induced by protons and alphas. In principle all the nuclei with $A > 11$ in interstellar space are target candidates. In practice, only ^{16}O, ^{12}C, and marginally ^{14}N are abundant enough to contribute appreciably. Alpha+alpha reactions were also important in the early days of the galaxy (Montmerle, 1977; Steigman and Walker, 1992).

Thanks mostly to the pioneering work of the Orsay group, the important excitation functions are now known with sufficient accuracy (reviewed in Reeves, 1974, and Read and Viola, 1984). They are displayed in Figs. 2, 3, and 11.

In the cosmic-ray sector, we would need to know the flux of H, ^{4}He, ^{12}C, ^{14}N, and ^{16}O as a function of energy throughout our galaxy (Figs. 4 and 5). Because of solar modulation effects, the low-energy parts of these fluxes are damped in the solar neighborhood. The problem of extrapolating outside the solar cavity is still not entirely solved. Data obtained by the *Voyager* satellite at 24 AU

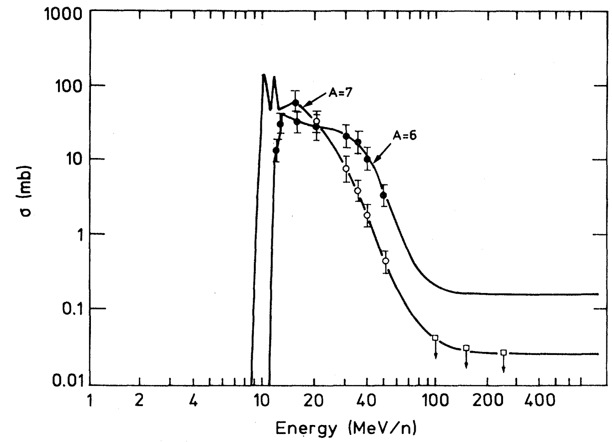

FIG. 11. The cross sections for the formation of ^{7}Li and ^{6}Li by $\alpha+\alpha$ reactions as a function of energy per nucleon. Note that the cross sections become very small above a few tens of MeV.

(farther than Saturn) have been used to obtain the best estimated fluxes and energy spectra (Ferrando *et al.*, 1991; Ferrando, 1992; Webber *et al.*, 1992). The omnidirectional flux of protons with energy larger than 100 MeV outside of the solar cavity is estimated to be 16 ± 4 cm^{-2} s^{-1}. Gamma-ray studies indicate that the flux decreases with galactocentric distance (Cesarsky *et al.*, 1977)

Each of the fast nuclei in the GCR process is faced with three possible fates. First, it may be destroyed by nuclear collision with interstellar atoms. The total destruction cross sections have been measured.

Second, it may also escape from the galaxy. The relevant parameter is the "escape length" Λ: the amount of matter met by a cosmic ray on its way from its source to the border of the galaxy. The value of this parameter is calculated from the abundance ratio of spallated nuclei (Li, Be, B) to the target nuclei (C, N, O); at the energies of interest in the present galactic medium $\Lambda \approx 6$ g cm^{-2}.

Third, it may be decelerated by collisions with interstellar electrons ("ionization losses") all the way down to interstellar thermal energies. (The energetics of the process has been analyzed by Ryter *et al.*, 1970, and constitutes yet another argument in favor of the GCR origin of the light elements as opposed to stellar origins.) Those are the atoms that will eventually be incorporated into stars and manifest themselves in the stellar spectra.

The number of thermalized nuclei added to the interstellar gas (dn_i/dt, in cm^{-3}s^{-1}) is the sum of two contributions:

(a) Spallation of interstellar heavy nuclei by fast protons and alphas. In this case, the recoil energy of the light nuclei is small (a few MeV per nucleon) and they suffer no further destruction or loss.

(b) Spallation of fast heavy nuclei by interstellar H and He. In this case we must calculate the "current" of particles in energy space $[-n_i(dE/dt)_i$, where $\Phi_i - n_i v_i$ (Φ_i is the flux, n_i the space density, and v_i the velocity], and compute its value at an appropriately low energy,

$$dn_i/dt = +\Sigma_j \int_0^\infty \{\sigma_{pji}(E')\Phi_p(E') + \sigma_{\alpha ji}(E')\Phi_\alpha(E')\} n_j dE' - n_i(E)(dE/dt)_i \ , \tag{21}$$

where j refers to the targets and $\sigma_{pji}(E')$ is the spallation cross section for the reaction $p + j \Longrightarrow i$.

The value of $\Phi_i(E)$ in space is related to the GCR source spectrum $q_i(E')$ and to the escape length out of the galaxy by

$$\Phi_i(E) = w_i(E)^{-1} \int_E^\infty dE' \left[\left[q_i(E') + \Sigma_i \int_0^\infty dE'' \Phi_j(E'')[(n_\alpha/n_p \sigma_{\alpha ji}) + \sigma_{pji}]/(M_p + M_\alpha n_\alpha/n_p) \right] \times \exp\{-[R_i(E') - R_i(E)]/\Lambda_i\} \right] , \tag{22}$$

where $R_i(E) = \int_0^\infty dE'/w_i(E')$ is the ionization range of particle i at energy E per nucleon (in g cm^{-2}); $w_i = [-(dE/dt)_i/\rho v_i]$. Here Λ_i is the loss range of energetic particles against the combined effects of (a) decay (if the nucleus is radioactive with period τ_i; γ_i is the relativistic time factor and ρ the density of the interstellar gas); (b) nuclear destruction, and (c) escape from the galaxy (escape length $= \Lambda_e$ in g cm^{-2}). Moreover σ_{pi} and $\sigma_{\alpha i}$ are the nuclear total destruction cross sections of nuclei by protons and alphas,

$$\Lambda_i^{-1} = (\rho v \tau_i \gamma_i)^{-1} + \{(\sigma_{pi} + n_\alpha \sigma_{\alpha i}/n_p)/(M_p + n_\alpha M_p/n_p)\} + \Lambda_e^{-1} \ . \tag{23}$$

B. Results of galactic cosmic-ray computations

To compare theory and observations, it is convenient to study separately two phases of galactic life: (1) the "recent" era of galactic life (the last ten billion years or so, corresponding more or less to Pop I stars) and (2) the early days of the galaxy (Pop II stars).

1. The recent galactic era

The fluxes and energy spectra of the GCR process have been extrapolated from the solar system taking into account the effect of solar modulation (Ormes and Prothero, 1983; Webber *et al.*, 1992). The extrapolated values become highly uncertain below one hundred MeV. Fortunately, most of the contribution to the production rate comes from energies where the fluxes are better known.

Meteoritic data on the production rate of nuclei by spallation reactions shows that the GCR fluxes have not varied significantly in the last few billion years (Lal and Peters, 1967; Zanda, 1988).

The interstellar target abundances of C and O are known to an accuracy better than a factor of two (Anders and Grevesse, 1989; Grevesse and Noels, 1993). We know also that they have remained approximately at their present values for the last ten billion years or so (Spite, 1992).

TABLE I. Production ratios in galactic cosmic-ray model.

^{6}Li/^{9}Be=5	^{7}Li/^{9}Be=7	Li/Be=12
^{10}B^{9}Be=5	^{11}B/^{9}Be=12	B/Be=17
^{7}Li/^{6}Li=1,4	^{11}B/^{10}B=2,5	

The ratios of the production rates for the various nuclei have been computed (Reeves and Meyer, 1978; Walker, Mathews, and Viola, 1985), normalized to the production rate of ^{9}Be. They are given in Table I. The uncertainties are less than a factor of two. The present abundance ratios of the light nuclei, discussed earlier, is given in Table II. The uncertainties in the element abundances represent a factor of two on each side.

Normalizing the time-integrated flux of GCR spallation to the abundance of Be, it appears that, within the uncertainties, the B/Be and the ^{6}Li/Be are satisfactorily accounted for. Extra sources of ^{11}B and ^{7}Li are required. The case of boron will be discussed presently. The problem of lithium will be presented later.

2. The problem of the boron isotopic ratio

The boron ratio observed in the solar system (^{11}B/^{10}B=4.05 compared to the value of ^{11}B/^{10}B$\approx$2.5 expected from high-energy GCR generation) could be related to the existence of fluxes of low-energy particles (tens of MeV) located somewhere in the galaxy, perhaps around the GCR acceleration sources (Meneguzzi and Reeves, 1975; Reeves and Meyer, 1978; Walker *et al.*, 1985). Because of the high rate at which these low-energy particles ionize the interstellar medium, these fluxes are unable to propagate throughout the galaxy. They are expected to remain confined close to their accelerating sources. Because of the high peak in the ^{14}N$(p,\alpha)^{11}$B cross section around 10 MeV, the formation of ^{11}B would be greatly favored in these regions.

One possible argument against this mechanism comes from the recent study of Prantzos *et al.* (1993a), which shows that there low-energy fluxes could have generated a too high Li/Be ratio if they existed in the early days of the galaxy. As we shall discuss later, in view of the low abundance of the C, N, and O species in those times, most of the Li came from the $\alpha+\alpha$ reaction at low energy.

At any rate, the possibility that this mechanism is the

TABLE II. "Recent" galactic abundances of Li, Be, and B.

Lithium/H $=2\times10^{-9}$
Boron/H $=2\times10^{-10}$
Beryllium/H $=1.3\times10^{-11}$
These ratios do not appear to have varied by more than a factor of two in the last five billion years or so.
^{7}Li/^{6}Li $=12.5$ (within five percent at solar birth; within twenty-five percent today).
Thus ^{6}Li/H $=2\times10^{-10}$
^{11}B/^{10}B $=4.05$ (within one percent at solar birth; unknown today).
Thus ^{10}B/H $=4\times10^{-11}$

answer to the boron isotopic ratio puzzle can be tested. The same low-energy protons would generate gamma rays from the excitation of ^{12}C and ^{16}O. In particular, the 4.4-MeV line from ^{12}C and ^{16}O and the 6.1- and 7.1-MeV lines from ^{16}O would necessarily be associated with these hypothetical fluxes of tens of MeV particles. Bloemen *et al.* (1993) have reported the detection of these lines in the Orion complex.

It is also of interest to consider other physical mechanisms that are known to produce appreciably higher ^{11}B/^{10}B ratios than nuclear spallation reactions. The photodisintegration, electrodisintegration, or neutrino disintegration of ^{12}C (Schaeffer, Reeves, and Orland, 1982; Boyd, Ferland, and Schramm, 1988; Woosley *et al.*, 1990) could in principle solve the boron isotopic problem if their time-integrated galactic ^{11}B contribution were closely equivalent (within less than a factor of two) to the GCR contribution. In view of the large uncertainties attached to their galactic yields, these possibilities can neither be confirmed nor ruled out. At any rate, the comparison between the GCR predicted and observed isotopic values can be used to put *an upper limit* on the nucleosynthetic importance of electro, photo-, and neutrino-disintegration processes in the galaxy: they should not have generated more ^{11}B than the galactic cosmic rays themselves!

Since we have no extra-solar-system data on the boron isotopic ratio, we should also consider the possibility of chemical or physical fractionation in the planetary system, with a preferential loss of ^{10}B. The fact that, as discussed previously, the meteoritic boron abundance ratio to hydrogen is some four times larger than the mean stellar value may be another indication of such processes. Stellar or galactic gas measurements of the boron isotopic ratio are needed to resolve this issue.

A possible solution in terms of a large neutron irradiation (Fowler *et al.*, 1962; the ^{10}B has a large neutron capture cross section) is essentially ruled out by the lack of a corresponding detectable isotopic effect on gadolinium (^{154}Gd has a much larger neutron capture cross section than the other Gd isotopes).

3. The early galactic era

The study of the GCR-generated elements in Pop II stars may yield important clues to the physical conditions accompanying the formation of galaxies (Ryan *et al.*, 1990, 1991; Rebolo *et al.*, 1988, 1993; Spite *et al.*, 1992; Duncan *et al.*, 1992; Gilmore *et al.*, 1991; Smith, Lambert, and Nissen, 1992). In Fig. 7 the abundances of Li, Be, and B are shown as a function of metallicity (Fe/H).

While the Li abundance reaches its big-bang plateau around one-tenth of the solar metallicity, the Be and B abundances are still going down with decreasing metallicity. At one part in one thousand of the solar metallicity, they have fallen by a factor of approximately one hundred below their present Pop I stellar value. This is qualitatively as expected from their GCR origin. This also can be used to give an upper limit to their hypothetical big-bang contribution, as is to be expected in the case of a strongly inhomogeneous big bang (no plateau similar to the lithium plateau has yet been found).

To obtain quantitative estimates of the GCR expected yields, we cannot simply use the standard GCR models developed for Pop I stars. We have to consider several factors that differentiated the ancient situation from the present one. These factors include the target abundances, the flux intensity of the galactic cosmic rays, and the manner of their propagation in the galaxy. The diffusion-like propagation involves, in turn, the matter density, the magnetic configuration governing the mean free path of the fast particles, and the propagation volume in relation to the changing shape of the galaxy (Prantzos *et al.*, 1993a, 1993b).

Consider first the evolution of the target abundances. As we move toward the past, the ratio CNO/H decreases progressively, while the α/H ratio remains almost constant. Since the C, N, and O group reactions generate the three elements Li, Be, and B, while the $\alpha+\alpha$ reactions generate only ^{7}Li and ^{6}Li, one expects the relative ratio of Li/BeB to increase with decreasing metallicity (Montmerle, 1977; Steigman and Walker, 1992). In this respect we note that, given the shape of the cross sections (Figs. 2, 3, and 11) and the shape of the GCR energy spectrum (Figs. 4 and 5), the $p+$CNO products are mostly created in the hundreds of MeV range, but the $\alpha+\alpha$ products are almost entirely generated in the tens of MeV range.

Consider next the evolution of the paleo-GCR flux spectrum both at its source and in space. At the source, we expect the injected power, presumably related to acceleration mechanisms in a supernova, to have been larger in the past, in relation to the expected higher star formation rate in the early days of the galaxy. This fact is best studied with galactic evolution models aimed at accounting for the variation of abundant elements (Fe,O,C) with galactic age or the number of stars of given metallicity as a function of metallicity (Pagel, 1987, 1989, 1992; Brown, 1992) and also of the relative abundance of long-lived radioisotopes (Reeves and Johns, 1976; Meyer and Schramm, 1986; Reeves, 1991; Colin and Schramm, 1992).

The numerical value of the escape length Λ of the galactic cosmic rays plays an important role in the formation of nuclei through the GCR bombardment of interstellar gas. Its present value ($\approx 6\ \mathrm{g\,cm^{-2}}$) governs the abundance of secondary nuclei (Li, Be, B, and also nuclei with mass number between Fe and Si) in the cosmic rays themselves. Since it is related to the shape of the galaxy, to the matter density, and to the configuration of the lines of force responsible for the diffusion of the fast particles, it may well have been different in the past.

The ratio of tens of MeV to hundreds of MeV particles in the interstellar flux of galactic cosmic rays is a key parameter of this discussion. In early times, the low value of the C, N, and O target abundances resulted in an increasingly important relative contribution of the $\alpha+\alpha$ reaction in generating Li (but not Be and B). This relative contribution is modulated by the ratio of the fluxes in the tens of MeV (for $\alpha+\alpha$) and hundreds of MeV (for $p+$CNO), a ratio that decreases with increasing Λ as more and more source particles are heavily decelerated before they have a chance to escape from the galaxy. The main point here is the identification of Λ as the phenomenological parameter most appropriate to integrate the landscape of the early galaxy in discussing the GCR product abundances.

The ratio of B to Be is rather insensitive to the value of the escape length. It never becomes smaller than 10, even for the largest escape length. Its value obtained in Pop II stars (10 ± 5) is compatible with the computed production rate of the GCR model, showing that both species are mostly of GCR origin, even in the very early days of the galaxy (Walker *et al.*, 1992; Thomas *et al.*, 1992).

The detection of ^{6}Li in one metal-poor star (Smith, Lambert, and Nissen, 1992) coupled with the upper limit of $^9\mathrm{Be/H} < 1.4\times10^{-13}$ corresponds to a ratio of $^6\mathrm{Li}/^9\mathrm{Be} > 50$. The discrepancy between this value and the ratio expected from the CNO contribution, $^6\mathrm{Li}/^9\mathrm{Be}(p+\mathrm{CNO}) \approx 5$, provides the first evidence for the contribution of $\alpha+\alpha$ reactions to the abundances of Li. It also shows that the GCR process did not contribute more than a few percent to the Pop II lithium abundance. Finally this ratio has been interpreted as a strong argument against a rotationally induced strong depletion of Li in the surface layers of Pop II stars, as this process would have depleted ^{6}Li by an unacceptable factor.

Such computations show how the abundances of Li, Be, and B can provide information on the physics of the early galaxy.

VII. INDIVIDUAL HISTORY OF THE COSMOLOGICAL NUCLEI

A. Deuterium

Deuterium is the first nucleus to appear in big-bang nucleosynthesis. Its abundance gave the first indication that the universe is not bound by nucleonic matter (Reeves, 1971; Geiss and Reeves, 1972). Deuterium is not generated by normal stellar synthesis. The only other source of D is the spallation of ^{4}He by galactic cosmic rays, which contributes only one part in one thousand of its present abundance.

During galactic evolution D is systematically destroyed by astration. Its abundance should decrease with time, as the difference between the protosolar value and the present value would seem to confirm. We have no data on D for objects older than the protosolar nebula. The highest lower limit on the baryonic density comes from its value in the distant past. Measurements of deuterium abundance in galaxies of low metallicities would be of prime importance in improving this estimate.

B. Helium-3

The helium-3 nucleus is produced in significant quantity both by BBN and by stellar synthesis, as an intermediate step in hydrogen burning. It is found primarily in the outer stellar layers, where the temperature is not high enough to complete helium burning all the way to ^{4}He. As a star moves toward the red giant branch, a fraction of these nuclei is burned into ^{4}He while another fraction is convected to the surface and ejected by stellar winds. Other processes in novae or other advanced stages of stellar evolution may further generate significant amounts of ^{3}He. As for D, the GCR contribution is insignificant.

The large scatter in the observed abundances of ^{3}He makes it difficult to identify the various contributions to its galactic abundance. We do not even know if, after big-bang nucleosynthesis, the galactic abundance of ^{3}He increases or decreases with time.

Satellite experiments have been set up to measure the $^3\mathrm{He}/^4\mathrm{He}$ ratio in the interstellar gas drifting into the solar system (Geiss, 1993). Lemoine *et al.* (1993b) have considered a method based on absorption lines against spectra of target stars.

C. Helium-4

After BBN production at a helium mass fraction level of $Y \approx 0.23$, the main effect of all the main-sequence stars has been to increase the galactic value to $Y \approx 0.30$, implying a transformation of $\approx 7\%$ of the hydrogen into ^{4}He, corresponding to an energy release of ≈ 0.5 MeV per nucleon.

Comparing this value to the present galactic luminosity brings us back to a milder version of the old Hoyle and

Tayler (1964) model. The standard unit for this discussion is the luminosity per unit mass of the sun: 2 erg/gm/sec or 2×10^{-12} eV/nucleon/sec. Spiral galaxies typically have a luminosity per unit mass of only one-tenth of this solar value. Assuming this luminosity to have been constant over the last 15 billion years, we find that this represents an integrated energy emission of $2\times10^{-13}\times1.5\times3\times10^{17}=9\times10^{4}$ eV, equivalent to the fusion or only one percent of H into He. A helium increase of 7% requires the mean galactic luminosity to have been at least five times larger than the present luminosity.

From the helium increase of 7% we may also compute the radiation energy density release per nucleon in the cosmos. With the baryonic density obtained from big-bang nucleosynthesis (ρ_b between 1 and 5×10^{-31} g cm^{-3}), the radiation density released is between 0.03 and 0.15 eV/nucleon, quite comparable to the density of the 3 K fossil radiation. In comparison, the stellar radiation density in the galactic plane is 0.4 eV/nucleon (Allen, 1973).

Because of universal expansion, the fraction of this radiation coming from intergalactic space is a function of the cosmological model adopted as well as of the chronology of helium burning. The interesting conclusion, however, is that a non-negligible fraction of the radiation energy density in the galactic plane comes from helium burned in distant galaxies.

D. Lithium-6

The lithium-6 nucleus is a pure product of GCR reactions. It comes mostly from the spallation of interstellar ^{16}O and ^{12}C by GCR protons. The report by Smith, Lambert, and Nissen (1992) of the detection of ^{6}Li/^{7}Li$\approx$0.05 in an old Pop II star is the first observational evidence of the $\alpha+\alpha$ contribution. Detection of this nucleus in stellar surfaces is an important tool for the study of the physics of stellar outer layers.

E. Lithium-7

Lithium-7 has the unique characteristic of owing its abundance to three different mechanisms, each contributing amounts which differ by less than one order of magnitude. The BBN contribution dominates the galactic gas abundance in the first billions of years. When the galactic mass fraction of heavy elements became larger than one part per thousand (one-tenth of the present value), a still unidentified stellar source managed to increase the abundance by an extra factor of ten.

The observations of lithium abundances in evolved stars of the asymptotic giant branch (Scalo, 1986) at a value appreciably larger than the Pop I value of 10^{-9} suggest that these stars play a role in this phenomenon. Other candidates such as novae and supernovae have also been suggested. Convincing quantitative models of the galactic enrichment are still lacking. The GCR production of ^{7}Li, evaluated through the abundance of ^{9}Be, never dominates the abundance curve.

F. Beryllium-9

Another pure product of GCR nucleosynthesis, Be is generated only by the bombardment of ^{12}C and ^{16}O. Its unique origin makes it a particularly useful monitor of time-integrated factors of galactic evolution such as the product of particle fluxes and abundance of targets. An early big-bang contribution, through the hypothetical effect of the quark-hadron phase transition, appears highly unlikely.

G. Boron-10

The formation mode of boron is the same as that for ^{9}Be. Only in the solar system have we been able to identify separately both isotopes of boron. Isotopic boron ratio measurements in stars would be of great value, especially since, contrary to the case of the lithium isotopic ratio, we do not expect differential depletion on the main sequence (except at the very cold end).

H. Boron-11

The ratio B/Be$\approx$10 in old stars can be taken as an indication of the fact that, taking into account the ratios of spallation cross sections, the isotope observed is mostly ^{11}B. The difference between the standard GCR prediction, ^{11}B/^{10}B$\approx$2.5, and the observation, ^{11}B/^{10}B$=$4, points out the need for another source of ^{11}B. A strong flux of MeV protons could do the job. There is also the possibility of the effects of energetic photons, electrons, or neutrinos on ^{12}C. And one cannot yet exclude the possibility of isotopic fractionation in the early solar nebula. This hypothesis is further indicated by the overabundance of boron in meteorites (Cl chondrites) compared with the stellar mean value.

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