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[1]U. Becker *et al.*, Nucl. Instrum. Methods 128, 593 (1975).

[2]D. Antreasyan *et al.*, in Proceedings of the 19th International Conference on High Energy Physics, Tokyo, Japan, August 1978 (to be published).

[3]P. D. Luckey *et al.*, in *Proceedings of the International Symposium on Electron and Photon Interactions at High Energies, Hamburg, 1965* (Springer, Berlin, 1965), Vol. II, p. 423.

[4]S. A. Berends *et al.*, Phys. Lett. 63B, 432 (1976).

[5]R. D. Field and R. P. Feynman, Nucl. Phys. B136, 1 (1978).

[6]H. G. Sanders, Diplomarbeit, Physikalischen Institut, Aachen, Report No. HEP 74/07, 1974 (unpublished).

[7]T. A. Gabriel and B. L. Bishop, Nucl. Instrum. Methods 155, 81 (1978), and the references therein.

[8]H. Terazawa, Rev. Mod. Phys. 45, 615 (1973).

Cosmological Constraints on New Stable Hadrons

C. B. Dover[(a)]
Department of Physics, University of Pennsylvania, Philadelphia, Pennsylvania 19104

and

T. K. Gaisser and Gary Steigman
Bartol Research Foundation of The Franklin Institute, University of Delaware, Newark, Delaware 19711
(Received 16 February 1979)

Possible new hadrons containing massive stable quarks (e.g., color-sextet quarks) surviving as relics of the early stages of the big bang should be present in $Z>1$ nuclei at levels accessible to experiment ($\sim 10^{-10}$). Grand unified theories that explain the observed baryon asymmetry of the universe should not contain these new stable quarks unless they are prevented from evolving asymmetrically. Otherwise, the new hadrons would be as common as nucleons, which is clearly not the case.

Existing bounds on anomalously heavy stable isotopes of hydrogen are exceedingly low, $\sim 10^{-18}$ relative to ordinary hydrogen for masses less than 16 GeV.[1] In contrast, bounds on anomalous stable isotopes of nuclei with $Z>1$ are much less severe.[2] A new heavy quark prevented from decaying by a new conserved quantum number would, along with the usual light quarks, form a stable heavy hadron. In terrestrial material, such heavy hadrons may reside preferentially in $Z>1$ nuclei. There is, moreover, no reason to expect that a nucleus containing such a new, heavy object would have an integral atomic mass.

The limit on anomalous nuclei of arbitrary mass comes simply from comparisons of the masses of elements determined chemically with those determined physically by averaging the known isotopes.[2] In contrast, techniques to detect specific stable isotopes[3] can have sensitivities as great as one nucleus in 10^{16}. Experiments designed specifically to search for anomalous heavy nuclei should therefore easily eliminate the new heavy hadrons—or find them if they exist.

The possibility that there exist new, heavy, stable quarks (for our purposes the lifetime should exceed the age of the Universe $\sim 10^{10}$ yr) has been suggested recently by several authors.[4-7] The negative results of searches at accelerators[8] suggest that it is unlikely that the ~5-GeV b quark (constituent of the upsilon)[9] is stable as was originally suggested by Cahn.[5] Indeed, these experiments eliminate long-lived ($\tau \gtrsim 5\times10^{-8}$ sec) hadrons (production cross section $>\frac{1}{10}\sigma_T$) with mass $\lesssim$5–10 GeV. There is, however, no evidence at present against the existence of a stable hadron with mass $\gtrsim$10 GeV. There is, however, no evidence at present against the existence of a stable hadron with mass $\gtrsim$10 GeV. For example, the color-sextet quark of charge $+\frac{1}{3}$ that arises naturally in the theory of supergravity[6] based on SO(8) could be sufficiently massive to have escaped detection. A color-sextet quark would be unable to decay to the familiar (color-triplet) quarks, by the usual weak interaction, and the lowest-mass hadronic state H containing such a quark would be stable.

Quarks of high mass would have been produced copiously during the very early epochs of the big bang. Those heavy quarks which survive annihilation would subsequently have been confined in heavy hadrons which would annihilate further. The remaining heavy hadrons originating from the big bang should be present in terrestrial matter.

At early times when the density and temperature were sufficiently high, particle-antiparticle pairs were continuously being produced and annihilated. As the temperature decreases to $T < m$, it becomes (exponentially) difficult to create new pairs; annihilation may still continue. As long as equilibrium can be maintained the density of particles will have its equilibrium value[10]; for $T \ll m$,

$$n \approx n_{\text{eq}} = \frac{g}{(2\pi)^{3/2}} \left(\frac{kT}{\hbar C}\right)^3 \left(\frac{mc^2}{kT}\right)^{3/2} \exp\left(-\frac{mc^2}{kT}\right), \quad (1)$$

where g is the statistical weight. The annihilation rate $\Gamma = n_{\text{eq}}\langle \sigma v\rangle$, where $\langle \sigma v\rangle$ is the thermally averaged annihilation rate coefficient, then decreases rapidly as T decreases further. Very quickly, a critical temperature is reached such that the annihilation rate Γ and the expansion rate t^{-1} are comparable:

$$\Gamma(T_*) t(T_*) \approx 1. \quad (2)$$

Subsequently ($T < T_*$), production can be neglected relative to annihilation and the density evolves as

$$d(nV)/dt = -n^2 \langle \sigma v\rangle V. \quad (3)$$

In Eq. (3), V is the (time-dependent) volume of an arbitrary comoving volume.[11] Equation (3) is then to be solved, subject to the boundary condition that $n(T^*) \approx n_{\text{eq}}(T_*)$.

A rough analytic result[11] for T_*, which agrees rather well with Wolfram's[12] numerical integration of the full rate equation, is

$$T_* \approx m/\ln(10^{34} m g \langle \sigma v\rangle). \quad (4)$$

In Eq. (4) the mass is in GeV and $\langle \sigma v\rangle$ is in units of $\text{cm}^3\ \text{sec}^{-1}$.

So far the analysis is quite general and is applicable to the survival of symmetric pairs of any leptons or hadrons. In particular, for nucleons[13] $\langle \sigma v\rangle \sim 10^{-15}\ \text{cm}^3/\text{sec}$ and $T_*^{N\bar{N}} \sim 25$ MeV. In contrast, for heavy point quarks of mass m, $\langle \sigma v\rangle \sim 10^{-17}\ [m(\text{GeV})]^{-2}$ and $T_* \gtrsim 200$ MeV for $m \gtrsim 8$ GeV. At such temperatures, the interparticle spacing[14] is about $\frac{1}{2}$ fermi and quarks are still unconfined. As the temperature drops further, confinement occurs (at $T_c \lesssim T_*$) and heavy quarks also are confined into hadrons H. Since the cross section of $H\bar{H}$ annihilation is presumably comparable to (or somewhat smaller than) that for $N\bar{N}$ annihilation (related to the bag size and quark counting), annihilation resumes until condition (2) is fulfilled with $\langle \sigma v\rangle \lesssim 10^{-15}\ \text{cm}^3/\text{sec}$.

We must therefore solve Eq. (3) with $\langle \sigma v\rangle = \langle \sigma v\rangle_0 \approx 10^{-17}\ [m(\text{GeV})]^{-2}$ for $T_c < T < T_*$ and with $\langle \sigma v\rangle = \langle \sigma v\rangle_1 \lesssim 10^{-15}\ \text{cm}^3/\text{sec}$ for $T < T_c$. The solution,[11] subject to the boundary condition $n_H(T_*) = n_{\text{eq}}(T_*)$, is

$$T^3/n_H(T) = T_*^3/n_H(T^*) + 2\langle \sigma v\rangle_0 A(T^* - T_c) + 2\langle \sigma v\rangle_1 A(T_c - T) \to 2\langle \sigma v\rangle_1 A T_c, \quad T \ll T_c. \quad (5)$$

This result is not substantially changed if $T_* \lesssim T_c$.[15] In contrast, for nucleons, with $T_*^{N\bar{N}} < T_c$, the relevant cross section is always $\langle \sigma v\rangle_1$ and the appropriate solution of Eq. (3) is $T^3/n_{N\bar{N}} \to 2\langle \sigma v\rangle_1 A T_*^{N\bar{N}}$.

Since photons are always relativistic, $n_\gamma \propto T^3$, and it is convenient to compare to n_γ. Numerically[14] $n_\gamma \sim 10^{40.5}\ [T(\text{GeV})]^3$ around $T = T_c$. Later as lighter pairs annihilate (e.g., e^+e^-, $\mu^+\mu^-$) extra photons are created; this increases the number of photons in every comoving volume by about a factor of 4.[16] Thus at present we expect $n_H/n_\gamma \gtrsim 2 \times 10^{-20}$ and $n_{N\bar{N}}/n_\gamma \sim 2 \times 10^{-19}$. In contrast, $n_{N\bar{N}}/n_\gamma \approx 10^{-(9\pm 1)}$ experimentally. This well-known, spectacular discrepancy clearly implies that the evolution of the Universe was far from symmetric. We return to this point and its implication for heavy, stable hadrons in the conclusion.

Continuing with the assumption of H-$\bar{H}$ symmetry, but taking the observed value of the nucleon-to-photon ratio, one has

$$n_H/n_N \sim 2 \times 10^{-11}. \quad (6)$$

Note that since $\langle \sigma v\rangle_1$ and T_c are independent of m_H, Eq. (6) is also independent of the mass of the heavy hadron.

To go further and estimate the abundance of the heavy, stable (symmetrically produced) hadrons to nucleons in *terrestrial* matter we must now follow their evolution through primordial nucleosynthesis and condensation into the galaxy, the solar system, and Earth. At this stage we consider specifically the color-sextet quark, h, suggested by supergravity.[6] It has charge $\frac{1}{3}$ and would form color-singlet states with two triplet antiquarks ($\overline{q}$) coupled into an antisextet. The lightest states would presumably involve u and d quarks. We use the bag model to find the lightest state,[17] which is stable against weak decays.

The color magnetic interaction which determines the mass splittings of hadrons is

$$H = -\sum_{ij} \vec{\sigma}_i \cdot \vec{\sigma}_j \lambda_i \cdot \lambda_j M_{ij},$$

where M_{ij} involve spatial properties of the bag and wave functions. For a three-quark system with one heavy quark the coupling between the light quarks is dominant. For the (antisymmetric) color $\underline{3}^*$ diquark $\lambda_1 \cdot \lambda_2 < 0$, whereas $\lambda_1 \cdot \lambda_2 > 0$ for a $\underline{6}_c^*$ (symmetric).[18] Thus, whereas for a color-triplet diquark coupled to a heavy (e.g., strange or charmed) quark the light diquark has spin 1 for the least-massive configuration, for a sextet diquark it has spin 0. Overall antisymmetry of the diquark system (isospin×spin×color) then implies that in both cases the least-massive state has $I=0$. Thus $m_\Lambda < m_\Sigma$, $m_{\Lambda_c} < m_{\Sigma_c}$, and the least-massive states with sextet quarks are $\overline{H}^0 = \overline{h}(ud)_{I=0}$ and $H^0 = h(\overline{u}\,\overline{d})_{I=0}$.

Both H^0 and $\overline{H}^0$ are expected to be bound in $Z>1$ nuclei. Note, however, that H^0 may annihilate in a nucleus via the reaction $H^0 + N \to M^+ + \pi$, if energetically allowed, where $M^+ = huG$ is the lowest-lying stable "meson" containing an h quark (G is a gluon). Simple quark-counting arguments suggest that the coupling constants of $\overline{H}^0$ to ordinary scalar or vector mesons ϵ and ω are $\frac{2}{3}$ times the corresponding meson-nucleon couplings.[15] This enables us to estimate the depth V_0 of the Hartree potential (as in Ref. 19) for $\overline{H}^0$ in a $Z>1$ nucleus; we obtain $V_0 \approx 20$–30 MeV. For a very massive $\overline{H}^0$, this will also be close to the binding energy of the 1s state, to which the $\overline{H}^0$ always migrates (in the absence of Pauli restrictions). The binding of H^0 in nuclei should be significantly stronger because of the attractive coupling of ω to light antiquarks. M^+ should also be bound, but presumably less strongly than $\overline{H}^0$.

The binding of h hadrons in $Z>1$ nuclei ensures that such hadrons would be protected by Coulomb repulsion from annihilation during Galactic and solar-system condensation. The $p\overline{H}^0$ system is not bound because single-pion exchange is absent and the second-order pion tensor force (important for deuterium) vanishes in the static limit.[15] Even if the other $Z=1$ states (e.g., pH^0 or $pn\overline{H}^0$) are bound, it is likely that they would to a large extent be processed in primordial nucleosynthesis into $Z>1$ nuclei. Moreover, any charged, symmetrically produced stable heavy hadrons not incorporated in $Z>1$ nuclei would probably have annihilated since condensation of the Galaxy and solar system.[15] It is therefore possible that new species of stable heavy hadrons with masses as low as ~10 GeV could be hidden in $Z>1$ nuclei.

Since ~90% of all nucleons are protons ($Z=1$),[20] the fractional abundance of heavy hadrons in $Z>1$ nuclei is $\sim 2\times 10^{-10}$ [compared with Eq. (6)]. If the heavy hadrons were uniformly distributed among all nuclei with $Z>1$ then, the abundance per mass A nucleus would be $2\times 10^{-10} A$.

As we emphasized earlier, this result is for a universe symmetric in H-$\overline{H}$. There is another important possibility. The observed nucleon excess may prove to be the consequence of the physics of baryon-number–nonconserving, CP-nonconserving, grand unified theories in the context of the very early ($T \gtrsim 10^{15}$ GeV) nonequilibrium stages of evolution of the Universe.[21] If this should prove to be correct, then it may be expected that the new heavy quarks would also evolve asymmetrically. In view of limits such as that of Ref. 2, it would appear that any new stable quarks in such theories will have to be introduced in such a way that the h-$\overline{h}$ asymmetry is reduced by at least several orders of magnitude relative to the N-$\overline{N}$ asymmetry.

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[a] On leave from Brookhaven National Laboratory, Upton, N. Y. 11973.

[1]Richard A. Muller, Luis W. Alvarez, William R. Holley, and Edward J. Stephenson, Science 196, 521 (1977), and references therein, especially T. Alvager and R. Naumann, Phys. Lett. 24B, 647 (1967).

[2]See, e.g., J. Schiffer, BNL Report No. BNL Pub. 50445, 1974 (unpublished), p. 61; R. J. Holt *et al.*, Phys. Rev. Lett. 36, 183 (1976); P. B. Price and J. Stevenson, Phys. Rev. Lett. 34, 409 (1975). Limits from mass spectrographic searches for specific isotopes (integer A) are at a similar level (Walter H. Johnson, Jr., private communication).

[3]See, for example, C. L. Bennett *et al.*, Science 201, 345 (1978). Sensitivities can be increased by starting (as in Ref. 1) with a sample enriched in heavy isotopes. See also Richard A. Muller, Phys. Today 31, No. 2, 23 (1979).

[4]F. Wilczek and A. Zee, Phys. Rev. D 16, 860 (1977); A. Zee, to be published; the sextet quark discussed here is a "shiny" quark in the language of these authors. See also the discussion of G. Karl, Phys. Rev. D 14, 2374 (1976).

[5]B. W. Lee and R. E. Shrock, Phys. Rev. D 17, 2410 (1978); R. Cahn, Phys. Rev. Lett. 40, 80 (1978).

[6]M. Gell-Mann, Bull. Am. Phys. Soc. 22, 541 (1977), and to be published; D. Z. Freedman, in *Particles and Fields—1977*, AIP Conference Proceedings No. 43, edited by P. A. Schreiner, G. H. Thomas, and A. B. Wicklund (American Institute of Physics, New York, 1978), p. 65.

[7]Other work on exotic quarks includes E. Ma, Phys. Lett. 58B, 442 (1978); H. Fritzsch, Phys. Lett. 78B, 611 (1978); Y. J. Ng and S.-H. H. Tye, Phys. Rev. Lett. 41, 6 (1978); S. Pakvasa and S. F. Tuan, to be published; P. G. O. Freund and C. T. Hill, to be published.

[8]D. Cutts *et al.*, Phys. Rev. Lett. 41, 363 (1978); R. Vidal *et al.*, Phys. Lett. 77B, 344 (1978).

[9]We assume the conventional interpretation of the $\Upsilon(9.6)$ of Lederman and co-workers as $\bar{b}b$ (cf. Vidal *et al.*, Ref. 8).

[10]Ya. B. Zeldovich, L. B. Okun, and S. B. Pikelner, Phys. Lett. 17, 164 (1965), and Usp. Fiz. Nauk. 87, 113 (1965) [Sov. Phys. Usp. 8, 702 (1966)].

[11]The scale factor a that characterizes a comoving volume $V \propto a^3$ in the early Universe evolves as $a \propto 1/T$. Combining the equation for gravitational expansion, $G\rho t^2 = \text{const}$, with the expression for the energy density of all relativistic particles gives $t = A/T^2$ with $A \sim 10^{-6}$ GeV^2 sec at very early times [see, e.g., G. Steigman, Cargèse Lect. Phys. 6, 505 (1973), and Annu. Rev. Astron. Astrophys. 14, 339 (1976)].

[12]S. Wolfram, to be published.

[13]The $N\bar{N}$ annihilation cross section for $p_{\text{lab}} = 0.5\text{–}2$ GeV/c can be parametrized as $\sigma_A = 38 + 35/p_{\text{lab}}$ mb. See T. E. Kalogeropoulos, in Proceedings of the Fourth International Experimental Meson Spectroscopy Conference, Northeastern University, Boston, 1974 (unpublished).

[14]For a gas of relativistic particles $n_{\text{eq}} \approx (g/\pi^2)(kT/\hbar c)^3$, with $g \lesssim 100$ at $T > 200$ MeV (see, e.g., Steigman, Ref. 11).

[15]C. B. Dover, T. K. Gaisser, and G. Steigman, to be published.

[16]G. Steigman, K. Olive, and D. N. Schramm, to be published.

[17]T. deGrand *et al.*, Phys. Rev. D 12, 2060 (1975); we thank Professor J. Rosner for pointing out this argument to us.

[18]P. Carruthers, *Introduction to Unitary Symmetry* (Wiley-Interscience, New York, 1966), p. 74.

[19]C. B. Dover and S. Kahana, Phys. Rev. Lett. 39, 1506 (1977).

[20]D. N. Schramm and R. V. Wagoner, Annu. Rev. Nucl. Sci. 27, 37 (1977).

[21]M. Yoshimura, Phys. Rev. Lett. 41, 381 (1978), and 42, 746(E) (1979); S. Dimopoulos and L. Susskind, Phys. Rev. D 18, 4500 (1978), and Stanford University Report No. ITP-616 (to be published); D. Toussaint, S. Treiman, F. Wilczek, and A. Zee, Phys. Rev. D 19, 1036 (1979); J. Ellis, M. K. Gaillard, and D. V. Nanopoulos, CERN Report No. TH. 2596, 1978 (to be published), and Erratum (to be published); S. Weinberg, Phys. Rev. Lett. 42, 850 (1979).