Baryon-Symmetric Baryogenesis

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(Received 13 September 1989)

We propose a new scenario of baryogenesis in which baryon number is an exact symmetry of the fundamental Lagrangian. If the global U(1) symmetry associated with baryon number is spontaneously broken at early times (when the temperature was of order 1–100 GeV), then the out-of-equilibrium decay of baryonic scalar particles can yield a cosmic baryon asymmetry in ordinary matter. This baryon number is exactly compensated by antibaryon number in the vacuum. At low temperatures, the symmetry is restored, and the antibaryons show up as neutral (antibaryonic) scalar particles or as nontopological solitons.

PACS numbers: 98.80.Cq, 11.30.Qc, 14.80.Pb

The Earth, our solar system, and perhaps the entire observable Universe are composed of protons and electrons rather than antiprotons and positrons. This fundamental observation implies that early in the history of the Universe, there must have been a small excess in the number of protons as compared to the number of antiprotons; if not, then baryon-antibaryon annihilations would have completely wiped out both matter and antimatter. The baryon excess is usually characterized by the dimensionless ratio B, defined to be the difference between the number density of baryons and antibaryons divided by the entropy density. Today, $B \simeq 10^{-10}$. In 1967, Sakharov[1] set down the ingredients necessary to produce a baryon excess. The most basic observation is that if all fundamental interactions conserve baryon number, then the baryon number today depends on initial conditions.

The grand unified theories (GUT's) of the 1970s opened up the possibility that baryon number is not conserved. Generically, GUT's predict that the proton will decay at a rate $\sim g^4 m_p (m_p/m_X)^4$, where m_X is the mass of the boson which mediates baryon-violating interactions, g is the coupling constant for these interactions, and m_p is the mass of the proton. (Here and throughout, we use units in which $\hbar = c = k_B = 1$.) Today, the rates for proton decay and other baryon-violating interactions are highly suppressed because m_X is very large ($m_X \sim 10^{15}$ GeV). However, in the early Universe, when temperatures were of order m_X, baryon-violating interactions would have occurred at a rate comparable to other processes. So long as Sakharov's other conditions (C and CP violation and a departure from thermal equilibrium) were met, the proton excess could have been *generated* by fundamental process[2] and would therefore be independent of initial conditions.

In this Letter we propose a different paradigm for baryon violation; one which allows—perhaps even forces—the baryon excess to be generated at relatively low temperatures ($T \sim 1$–100 GeV). We assume that the fundamental Lagrangian governing all interactions is completely *baryon symmetric;* that is, it is invariant under the simultaneous [$U(1)_B$] transformation

$$\psi_i \rightarrow e^{ib_i\theta}\psi_i , \tag{1}$$

where ψ_i are the fields in the Lagrangian and b_i are their baryon numbers. However, we also assume that baryon number is *spontaneously* broken [i.e., that at least one field carrying baryon number gets a vacuum expectation value (VEV)] at early times. As long as Sakharov's other requirements are met, an excess of baryons over antibaryons can develop. Why then is baryon number apparently not violated today? As the Universe expands and cools, the VEV's of the baryonic fields disappear (or at least get very small); thus at *low* temperatures $U(1)_B$ is restored. Although low-temperature symmetry restoration is counterintuitive, it has been observed in the ferroelectric behavior of Rochelle salts[3] and has been analyzed for a variety of reasons by many authors.[4]

One consequence of our assumptions has no analog in scenarios with explicit baryon violation. A baryon-symmetric Lagrangian leads to equations of motion $dB/dt = 0$. When baryon number is spontaneously violated, this becomes[5]

$$\left(\frac{dB}{dt}\right)_{\text{particles}} + \left(\frac{dB}{dt}\right)_{\text{vacuum}} = 0 . \tag{2}$$

Therefore, any baryon-antibaryon excess generated is exactly compensated by an excess in the vacuum. For example, if the field which gets a nonzero VEV is a scalar field, then its baryonic charge density is $n_B|_{\text{vacuum}} = ib_\phi(\phi^*\dot{\phi} - \dot{\phi}^*\phi)$. The VEV of ϕ must therefore be nonzero and also time dependent in order to compensate any baryon asymmetry in particles. If the Universe starts out with $B = 0$, then Eq. (2) tells us that the baryon density in the vacuum is equal and opposite to the baryon density in particles: $n_B|_{\text{vacuum}} = -n_B|_{\text{particles}}$. At low temperatures, when the VEV's go away (or get small), this antibaryon number must show up somewhere. We will offer several suggestions as to what form this antimatter might take today.

The fact that the Universe contains as much antimatter as matter suggests that the energy scale for baryon-symmetric baryogenesis is relatively low ($\sim$1-100 GeV). This follows since the total energy density of the antimatter today—in whatever form it takes—is $\rho_{\text{AM}} = n_B(E/Q)$, where E/Q is the energy per baryon number of the antimatter. This energy density cannot exceed the critical density of the Universe: $\Omega_{\text{AM}} \equiv \rho_{\text{AM}}/\rho_c < 1$. Since protons contribute $\Omega_B \sim 10^{-1}$-10^{-2}, E/Q of antimatter must be less than $(10\text{-}100)m_p$. Naively, and indeed in most of the models we have examined, the interesting physics responsible for baryogenesis must be in this energy regime. A low-temperature scenario for baryogenesis has a number of advantages over high-temperature (e.g., GUT-inspired) models. For example, the reheating temperature in inflationary models must be above the scale of baryogenesis in order that the baryon asymmetry produced not be diluted during the inflationary epoch. GUT-scale baryogenesis therefore requires a high reheating temperature, and this is difficult to achieve in realistic models. On the other hand, baryon-symmetric baryogenesis is compatible with a much lower reheating temperature.

In what follows, we will present our main results. Detailed calculations will appear in a forthcoming publication.[6] As discussed above, $U(1)_B$ must be spontaneously broken at high temperatures. A simple "toy" model involving two complex scalar fields ϕ and σ illustrates how this might arise. Consider the Lagrangian $\mathcal{L} = |\partial_\mu \phi|^2 + |\partial_\mu \sigma|^2 - V(\phi,\sigma)$, where

$$V(\phi,\sigma) = m_\phi^2|\phi|^2 + \alpha_1|\phi|^4 + \alpha_2|\sigma|^4 - 2\alpha_3|\phi|^2|\sigma|^2, \quad (3)$$

and $m_\phi^2 > 0$ is the mass of ϕ. The α_i are assumed to be real and positive. Furthermore, $\alpha_1\alpha_2 > \alpha_3^2$ is required to ensure that the potential be bounded from below. $V(\phi,\sigma)$ has its minimum at $\langle\phi\rangle = \langle\sigma\rangle = 0$. By definition, this is the zero-temperature vacuum state. However, at finite temperatures there are quantum thermal corrections to V which change this state. Using standard techniques,[4] we calculate the effective (one-loop) temperature-dependent potential $V(\phi,\sigma;T) = V(\phi,\sigma) + V_1(T)$, where

$$V_1(T) = \frac{T^2|\phi|^2}{6}(2\alpha_1 - \alpha_3) + \frac{T^2|\sigma|^2}{6}(2\alpha_2 - \alpha_3). \quad (4)$$

If we assume that $2\alpha_1 < \alpha_3 < 2\alpha_2$, then when $T > T_c$ $[\equiv \sqrt{6}m_\phi/(\alpha_3 - 2\alpha_1)^{1/2}]$ the minimum of $V(\phi,\sigma;T)$ is at $\langle\sigma\rangle = 0$ and $\langle\phi\rangle = \kappa(T^2 - T_c^2)^{1/2}$, where $\kappa \equiv [(\alpha_3 - 2\alpha_1)/12\alpha_1]^{1/2}$. That is, at temperatures above the critical temperature T_c, the U(1) symmetry associated with rotations of ϕ is spontaneously broken. In passing, we note that σ need not be a single complex scalar and in fact could be the standard-model Higgs doublet.[6]

Now consider the interaction Lagrangian

$$\mathcal{L}_{\text{int}} = \lambda_2 m_\phi \phi^* \phi_1 \phi_2^* + \lambda_3 m_\phi \phi^* \phi_1 \phi_3^* + f_1 \phi_1 U^T C D + f_2 \phi_2^* U^T C E + f_3 \phi_3^* U^T C E + \text{H.c.}, \quad (5)$$

where $SU(3)_{\text{color}}$ indices have been suppressed and C is the charge-conjugation matrix. U, D, and E refer to ordinary quarks and leptons. For simplicity, we consider only one generation, and, in the particular, the heaviest generation so that U, D, and E are the top quark, bottom quark, and τ lepton, respectively. Furthermore, we assume that the fermions are right-handed, $SU(2)_L$ singlets. The generalization to left-handed fermions or to more than one generation is straightforward. The Lagrangian in Eq. (5) is symmetric under a global $U(1)_B$ symmetry [Eq. (1)]. The $SU(3)_{\text{color}} \times U(1)_{\text{em}} \times U(1)_B$ content of the scalar fields, as inferred from Eq. (5) and the known quantum numbers of the quarks and leptons, are $\phi(1,0,-1)$, $\phi_1(3,-\frac{1}{3},-\frac{2}{3})$, $\phi_2(3,-\frac{1}{3},\frac{1}{3})$, and $\phi_3(3,-\frac{1}{3},\frac{1}{3})$. We see that ϕ has baryon number -1, so when ϕ gets a VEV, $U(1)_B$ is spontaneously broken. For further reference we note that we could have assigned ϕ baryon number $-\frac{1}{2}$ so that the $U(1)_B$ invariant couplings become $\lambda_i \phi^* \phi^* \phi_1 \phi_i^*$. The development of a baryon asymmetry is similar in these two models, but the subsequent behavior of the Universe differs dramatically.

A nonzero VEV for ϕ implies that there are mass-mixing terms for ϕ_1, ϕ_2, and ϕ_3. These are represented in Feynman diagrams by propagators that change one ϕ_i to another as in the diagram for the process $\phi_1 \to UE$ displayed in Fig. 1. Since ϕ_1—and the same holds true for ϕ_2 and ϕ_3—couples to two modes ($\bar{U}\bar{D}$ and UE) with different baryon numbers ($-\frac{2}{3}$ and $\frac{1}{3}$) it can no longer be assigned a baryon number. Thus, when $U(1)_B$ is spontaneously broken, a baryon-symmetric Lagrangian becomes similar to ones considered in GUT-based models.[2]

Let us now turn to the early Universe, and specifically to a time when the temperature was above T_c. We assume that at some temperature T_D (less than the masses m_i of ϕ_i but greater than T_c) there are equal number densities n_0 of ϕ_1, ϕ_2, and ϕ_3 particles and antiparticles. If the ϕ_i's are in equilibrium, then when they decay no asymmetry will be produced in the quark fields. However, if the ϕ_i's decay when they are out of equilibrium, then an asymmetry in the quark fields can develop. Thus we assume[7] that $n_0 \simeq T_D^3$; that is, there are many more ϕ_i's than there would be in equilibrium (in equilibrium $n \propto e^{-m_i/T_D}$). The final baryon excess (in ordinary matter) is then given by the asymmetry produced per decay times the number density of decaying particles at T_D. First we calculate ϵ_i, the net baryon number produced in the decay of ϕ_i and its antiparticle,

$$\epsilon_i \equiv \frac{1}{\Gamma_i}[\tfrac{1}{3}\Gamma(\phi_i \to UE) - \tfrac{2}{3}\Gamma(\phi_i \to \bar{U}\bar{D}) - \tfrac{1}{3}\Gamma(\bar{\phi}_i \to \bar{U}\bar{E}) + \tfrac{2}{3}\Gamma(\bar{\phi}_i \to UD)] . \tag{6}$$

Here, Γ_i is the total decay width of ϕ_i. The leading contribution to ϵ_i comes from interference terms between tree-level and one-loop graphs. The result (say, for $i=1$) is

$$\epsilon_1 = \frac{\mathrm{Im}(f_2^* f_3 \lambda_2 \lambda_3^*)\kappa^2(T^2-T_c^2)m_\phi^2}{(m_1^2-m_2^2)(m_1^2-m_3^2)}\left[\mathcal{J}\left(\frac{m_2^2}{m_1^2}\right)-\mathcal{J}\left(\frac{m_3^2}{m_1^2}\right)\right], \tag{7}$$

where $\mathcal{J}(\eta) = -(4\pi)^{-1}[1-\eta\ln(1+1/\eta)]$. (Here, we have taken $m_i^2 \gg \lambda_a |\langle\phi\rangle| m_\phi$.) We note that the coupling constants must be complex for there to be a nonzero ϵ; that is, CP must be violated. Furthermore, does not depend on the phase of ϕ though it does vanish if $\langle\phi\rangle = 0$. (If it did depend on the phase of ϕ, there would be small domains of baryons and antibaryons.) Finally, note that all five terms in Eq. (5) are necessary to get a nonzero ϵ. To calculate B, the ratio of this excess to the entropy density, we must take into account the increase in entropy due to the extra particles produced in decays. For simplicity, we can take all the m_i's to be of order m; then all the ϵ_i's are comparable ($\epsilon_i \sim \epsilon$) and

$$B \simeq \left(\frac{n_0}{(\pi^2/30)g_* T_D^3}\right)\left(1+\frac{3n_0 m}{(\pi^2/30)g_* T_D^4}\right)^{-3/4}\epsilon , \tag{8}$$

where g_* (~ 100) is the effective number of degrees of freedom when $T = T_D$. Clearly, with a reasonable choice of parameters, B can be $\sim 10^{-10}$. We note that both inverse decays (e.g., $UE \to \phi_i$) and baryon-violating scattering processes (e.g., $UE \to \bar{U}\bar{D}$) can potentially wash out the asymmetry. This is avoided if the respective rates are less than the expansion rate at T_D. The exact constraints depend sensitively on T_D and the top-quark mass. However, as long as T_D and m_ϕ are smaller than m, the constraints can be easily satisfied.

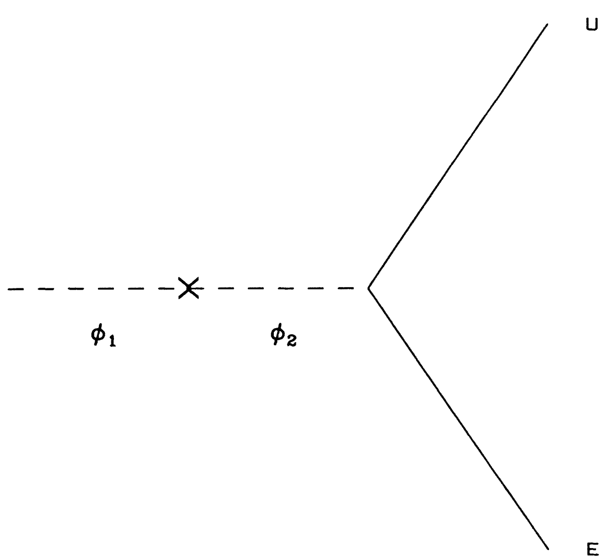

FIG. 1. Baryon-violating decay made possible by the vacuum expectation value of ϕ. The $\times$ represents the mass-mixing term $\lambda_2 m_\phi \langle\phi\rangle^* \phi_1 \phi_2^*$.

The baryon asymmetry generated in the quark fields is exactly compensated by a density of antibaryons in the vacuum. The simplest vacuum configuration with nonzero baryon density has $\phi = ve^{-i\omega t}/\sqrt{2}$ with $v^2\omega = n_B|_{\text{particles}}$. At temperatures above T_c, ϕ sits at the minimum of its potential [$v/\sqrt{2} = \kappa(T^2-T_c^2)^{1/2}$], and ω is just $n_B/2\kappa^{-2}(T^2-T_c^2)^{-1}$. When the temperature drops below T_c, the minimum of the ϕ potential is at $\phi = 0$. However, because the baryon number in the vacuum is nonzero, v cannot be zero everywhere. To gain some understanding as to how ϕ behaves today, let us assume that ϕ is homogeneous in space and consider the zero-temperature, small-oscillation expression for its energy density $\rho_\phi = (\omega^2 + m_\phi^2)v^2/2$. Substituting n_B/v^2 for ω and minimizing with respect to v we find that $v = (n_B/m_\phi)^{1/2}$, $\omega = m_\phi$, and $\rho_\phi = m_\phi n_B$. These results indicate that the ϕ field is behaving like a condensate of cold (zero-momentum) ϕ particles.

Observations today place several obvious constraints on the parameters of this theory. We know that m_ϕ must be greater than m_p or else the proton would decay into a ϕ^* and a positron. If the antibaryon number of the Universe today is in the form of cold ϕ particles, then an upper limit on m_ϕ follows from cosmological considerations. Assuming that most of the ϕ particles have not yet decayed, their energy density, as compared to the closure density ρ_c, is

$$\Omega_\phi \equiv \rho_\phi/\rho_c = 0.03h^{-2}(m_\phi/m_p)(B/10^{-10}) ,$$

where $H_0 = 100h$ km/sec Mpc is the Hubble parameter today. This places a tight upper bound on m_ϕ but also suggests the interesting possibility that ϕ particles close the Universe; that is, the dark matter is actually antimatter. An additional constraint comes from the fact that ϕ can decay into an antiproton and a positron, potentially wiping out the baryon asymmetry. We estimate the lifetime of the ϕ's to be $(1\ \text{sec})\lambda^{-2}f^{-4}(m_p/m_\phi)^9 \times [m/(1\ \text{TeV})]^8$. It is possible to make this lifetime greater than the age of the Universe by requiring the Yukawa couplings to light generations to be much smaller than those to the heavy generation, but this would require fine tuning the radiative corrections. A more natural model is one in which ϕ has baryon number $-\frac{1}{2}$. In this model, ϕ depletion takes place due to scattering $\phi\phi \rightarrow \bar{p}e^+$ and this rate is always much smaller than the expansion rate.

Until now we have assumed that once $U(1)_B$ was restored, the antibaryon number in the vacuum generated during the baryogenesis epoch was transferred to free ϕ particles. Actually, the lowest-energy configuration, given a fixed baryon number, may be quite different. We now discuss the intriguing possibility that the lowest-energy state is one with $\phi = 0$ everywhere except in certain localized regions of antibaryonic charge. That is, we suppose that the antimatter created during the baryogenesis epoch is found today in nontopological solitons (NTS's).[8] Indeed, following an epoch of baryon-symmetric baryogenesis, the Universe is ripe for NTS formation, since the vacuum necessarily contains a nonzero charge density (baryon number) associated with a global symmetry $[U(1)_B]$, precisely the situation postulated *ad hoc* in previous discussions of NTS formation.[9] The scenario is as follows: As T falls below T_c, ϕ's potential drives $\langle\phi\rangle$ to zero and $U(1)_B$ to symmetry restoration. However, the vacuum is charged with antibaryon number. If there is some attractive force between ϕ's, then NTS's (here, bubbles of antimatter or BAM's) will form. Before discussing specific models for BAM's, we mention that there is a model-independent upper limit on the charge of a BAM set by the charge within the horizon at the time when they form: $Q_{\text{max}} \simeq 10^{45}(B/10^{-10})(m_p/T_f)^3$, where T_f is the temperature when the BAM's form.

BAM's similar to Coleman's Q-balls[10] can form if the potential for ϕ is of the form $V(\phi) = m_\phi^2|\phi|^2 - \lambda|\phi|^4 + |\phi|^6/M^2$, where M is some mass scale and $\lambda > 0$. Of course, $V(\phi)$ must be an effective potential calculated from a more fundamental theory (presumably at the energy scale M) since ϕ^6 interactions are nonrenormalizable. Here, it is the attractive ϕ^4 interaction that stabilizes the Q-ball. The mass of the Q-ball is $E = Qm_\phi[1-(\lambda M/2m_\phi)^2]^{1/2}$ (clearly less than Qm_ϕ as it must be to ensure stability against dispersion into free particles) and its radius is $R \simeq (Q/\lambda M^2 m_\phi)^{1/3}$. A BAM with $Q = 10^{42}$ would be about a millimeter in size and weigh about 10^{19} g, roughly equal to the mass of a large mountain. BAM's can also form in a renormalizable theory with two or more scalar fields.[11]

Baryon-symmetric baryogenesis appears to be an attractive paradigm for explaining the observed proton-antiproton asymmetry. One definite prediction, that the number density of antibaryons be the same as the number density of baryons, distinguishes this scenario from scenarios in which baryon number is violated in the fundamental Lagrangian, and should lead to astrophysical consequences. Moreover, since the natural energy scale for baryon-symmetric baryogenesis is in the GeV range, signals should appear in particle-accelerator experiments.

It is a pleasure to thank C. Hill for stimulating conversations. We would also like to thank J. Bernstein, S. Carroll, G. Feinberg, J. Frieman, J. Jubas, and T. Pyne for helpful comments. This work began during a visit to the Aspen Center of Physics and we thank that institution for hospitality and support. This work was supported in part by the National Science Foundation (Grant No. PHY-8604396) and the U.S. Department of Energy (Grant No. DE-FG02-84ER40158 with Harvard University).

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