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Are antibaryons a signal for a phase transition in ultrarelativistic nucleus-nucleus collisions?

T. A. DeGrand
Department of Physics, University of Colorado, Boulder, Colorado 80309
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In this paper, it is speculated that antibaryon production in relativistic nucleus-nucleus collisions is a signal for a phase transition from a high-temperature Debye-screened and chirally symmetric phase to a low-temperature confined and chirally broken phase. In a chiral model for the transition where baryons are topological excitations the production of baryons and antibaryons takes place by the Kibble mechanism. A qualitative discussion of possible experimental signatures is given.

There has been considerable interest recently in the properties of hadronic matter at extreme temperature and pressure.[1] There is fairly good theoretical evidence that matter made of gluons, described by the pure gauge sector of QCD, undergoes a first-order phase transition at a critical temperature of several hundred MeV from a low-temperature confining phase to a high-temperature Debye-screened phase. The most recent Monte Carlo studies appear to suggest that matter made of quarks and gluons also undergoes a transition (which is probably first order) from a low-temperature confined phase in which chiral symmetry is broken to a high-temperature Debye-screened phase in which chiral symmetry is restored.[2]

Although theorists have presented numerous signals for the existence of a quark-gluon plasma or of the phase transition,[1] none of them are truly striking or easy to see. The reason for this is that energy scales involved in the transition are typically only a few hundred MeV, and thus most signals are difficult to distinguish from the ordinary soft hadronization associated with conventional hadron production.

Baryons or antibaryons may tell a different story. At c.m. energies of a few tens of GeV/nucleon the $\bar{p}/\pi^-$ ratio in the central region is around a few percent for pp collisions. If a quark-gluon plasma is formed with a temperature $T_c \approx 200$ MeV, then the equilibrium $\bar{p}/\pi^-$ ratio produced when the plasma converted to hadrons would be

$$\bar{p}/\pi^- \approx 2\exp(m_\pi - m_p)/T = 0.04 \ .$$

Thus, antibaryons are a good signal for interesting physics in heavy-ion collisions because they are not produced thermally or with conventional processes: any $\bar{p}/\pi^-$ signal greater than a few percent or showing unusual properties, to be described below, must be due to nonequilibrium processes, perhaps related to the existence of the phase transition. Conversely, baryon production in the central region is a poor signal because the nuclear fragmentation regions are not expected to be well separated in rapidity.

An example of interesting effects in the baryon sector due to the confinement-deconfinement transition has been given by Witten.[3] He is concerned with the transition in the early Universe; I will first give a variation of his argument as applied to heavy-ion collisions, and then a second argument where baryon production is due to the chiral-symmetry-breaking transition. These arguments do not guarantee experimentalists a large number of antibaryons (hereafter $\bar{B}$'s), but suggest interesting signals even if there is no anomalous $\bar{B}$ production.

A first-order transition from a high-temperature phase A to a low-temperature phase B proceeds (as T falls below T_c) by nucleating bubbles of phase B which expand to fill the system. When the transition is almost complete, the system consists of small (shrinking) bubbles of phase A surrounded by phase B. Now Witten reminds us that in the high-temperature phase of QCD baryon number is carried by light quarks, whereas in the $T=0$ world, baryon number is carried by heavy protons. As bubbles of the confined world are nucleated out of the quark-gluon plasma, it is energetically favored for any quarks or antiquarks which cannot be bound into mesons to be pushed back into the plasma phase. In the latter stage of the transition, the regions of plasma phase which have become enriched in baryon number are disconnected from one another. At that point several things can happen. Witten has discussed the scenario in which these regions remain quark droplets even at $T=0$, a completely new stage of matter. A more conventional scenario is that these regions condense into ordinary hadronic matter. Since the quark distributions on the surfaces of the colliding hadron bubbles were random, the number of quarks minus antiquarks will fluctuate from bubble to bubble, and bubbles with $n_q - n_{\bar{q}} = \pm 3, \pm 6, \ldots$ hadronize with a nonzero baryon number. One expects that the number of baryons plus antibaryons, $N_B + N_{\bar{B}}$, will be proportional to the number of bubbles originally nucleated, N_b. (Of course, $N_B = N_{\bar{B}}$.) Thus we have produced a nonthermal baryon density as a result of a first-order phase transition. Deferring a discussion of its fate for now, we turn to a second model of $B\bar{B}$ production in the transition.

Chiral models provide another plausible scenario for the transition. Since the critical temperature is no more than a few times the pion mass, it might be appropriate to describe the transition using not the full QCD theory but an effective Lagrangian which describes QCD at low energies. This effective Lagrangian contains an $SU(N_f) \times SU(N_f)$ multiplet of (N_f) fundamental fields ϕ which represent the pions $(N_f = 2)$ or π's, K's, η, and η' $(N_f = 3)$: The most general such model, given by Pisarski and Wilczek,[4] is

$$L = -\tfrac{1}{2}\mathrm{Tr}\partial_\mu\phi^\dagger\partial_\mu\phi - \frac{\mu^2}{2}\mathrm{Tr}\phi^\dagger\phi - \frac{g_1}{4!}\mathrm{Tr}(\phi^\dagger\phi)^2 - \frac{g_2}{4!}(\mathrm{Tr}\phi^\dagger\phi)^2 + G(\det\phi + \det\phi^\dagger) \ . \qquad (1)$$

Here G is a term which includes instanton effects. This model has been discussed by Pisarski and Wilczek using weak-coupling perturbation theory; they show that it has a first-order transition from a low-temperature chirally broken phase to a high-temperature chirally symmetric phase, for $N_F \geq 3$.

In addition to mesons, the model of Eq. (1) possesses topological excitations with baryonic quantum numbers, which however, shrink to zero size without the existence of extra terms. The appropriate extra term was provided by Witten:[5]

$$L_B = L + \frac{2}{15\pi^2 f_\pi{}^2}\epsilon^{\mu\nu\alpha\beta}\,\mathrm{Tr}\phi\partial_\mu\phi\partial_\nu\phi\partial_\alpha\phi\partial_\beta\phi \quad , \tag{2}$$

leading to an effective Lagrangian for mesons and baryons in which the baryons are topoloigical excitation of the ϕ fields.

The critical properties of L_B are identical to those of L because the extra soliton term is an irrelevant operator (in the renormalization-group sense), so the analysis of Pisarski and Wilczek carries through unchanged. It is now easy to see that baryon production in the chiral transition proceeds exactly like monopole production in the early Universe, through the Kibble mechanism.[6] In the high-temperature phase, the vacuum expectation value of ϕ is randomly oriented in space, so that its average $\langle\phi\rangle$ is zero. If the chiral transition is first order, then as the temperature falls it will proceed through the nucleation of bubbles of the true ground state, in which $\langle\phi\rangle \neq 0$. These bubbles will expand explosively, filling the plasma. $\langle\phi\rangle$'s in different bubbles need not be correlated, so that when bubbles coalesce to fill the space, it will be impossible for $\langle\phi\rangle$'s from different bubbles to align uniformly. One expects to produce topological knots: These knots are baryons and antibaryons. Naively, one expects the number of knots so produced to be comparable in size to the number of bubbles nucleated in the transition.

It is important to note that in this model baryon production does not depend on the order of the transition. It is a topological effect. If the transition were second order, bubble nucleation would not occur. The parameter which is related to number of produced baryons in a first-order transition, the bubble radius $R(T)$, would be replaced by the correlation length $\xi(T)$. At $T = T_c$, ξ becomes infinite. As described by Einhorn, Stein, and Toussaint,[7] ξ is also the shortest wavelength at which fluctuations in ϕ become unstable. If T varies with time, then at some T, $d\xi/dt > c$. But one expects that the actual instantaneous correlation length ξ, the length scale over which the field fluctuates in some direction, cannot increase faster than c, so when ξ reaches the length for which fluctuations are unstable, the symmetry breaks and topological defects are well defined. For B and $\bar{B}$ production to occur in a heavy-ion collision with a second-order transition, we need ξ less than the transverse size of the nucleus when this occurs.

Let me now sketch a simple picture of bubble formation during a first-order transition. I will approximate the equation of state for hadronic matter with a simple bag equation of state, appropriate to a first-order transition.[8] The parameters are B = bag constant and g_i (g_f), the number of relativistic degrees of freedom above (below) the transition: $g_i \approx 37$ for eight gluons and two flavors of quarks; $g_f = 3$ for three pions. The equations of state are

$$\epsilon_i(T) = \frac{\pi^2}{90} g_i T^4 + B, \quad P_i(T) = \tfrac{1}{3} g_i \frac{\pi^2}{90} T^4 - \tfrac{4}{3} B \quad , \tag{3}$$

$$\epsilon_f = 3P_f = \frac{\pi^2}{90} g_f T^4 \quad . \tag{4}$$

The critical temperature is defined by $P_i(T_c) = P_f(T_c)$ or

$$(g_i - g_f)\frac{\pi^2}{90} T_c{}^4 = 4B \quad ;$$

the latent heat is $4B$.

A simple model for the thermal nucleation for bubbles is obtained by assuming a surface tension S between domains of the different phases.[9] The critical radius of a bubble nucleated at a temperature T is

$$R_0(T) = \frac{2S}{\Delta P} \equiv \frac{R_0}{1 - \hat{T}^4} \tag{5}$$

and the nucleation rate

$$p(t) = P_0 T_c{}^4 \exp\left(-\frac{W_0}{\hat{T}(1-\hat{T}^4)^2}\right) \tag{6}$$

with

$$\hat{T} = T/T_c, \quad W_0 = \tfrac{2}{3} B V_0(1)/T_c \quad ,$$

where $R_0 = 3S/2B$ is a characteristic radius and

$$V_0(\hat{T}) = 4\pi R_0(T)^3/3 \quad .$$

Now we compute bubble-formation times in the central region of a relativistic heavy-ion collision. The calculation is similar to those for the early Universe,[10] with the important exception that since all distances are the same order of magnitude, the original volume at which the bubble was nucleated is an important contribution to the effective bubble volume. The fraction of space containing no bubbles is

$$f(t) = \exp\left[-\frac{4\pi}{3}\int_0^t dt' p(T(t'))[R_0(T)^2 + v^2(t-t')^2]^{3/2}\right] , \tag{7}$$

where v is the bubble's velocity, and I assume the space itself is not expanding. In the central rapidity region the solutions to the hydrodynamic equations for the expansion were given first by Landau.[11] One finds that the flow velocity $\theta(y,t)$ is equal to the rapidity y. If we assume that bubble nucleation begins at a time τ after the collision, the volume of plasma at a time $t+\tau$ is $A(t)(t+\tau)y$, where $A(t)$ is the cross-sectional area of the plasma. Then the number of bubbles (b) per unit rapidity is

$$\frac{dN_b}{dy} = \int_0^\infty dt\, p(T(t)) f(t) A(t)(t+\tau) \quad . \tag{8}$$

In order to integrate (7) and (8), one needs a relation between the time t and the temperature T of the plasma, as well as numerical values for v and R_0. The first relation may be obtained from the hydrodynamic calculations of Kajantie, Raitio, and Ruuskanen.[8] Allowed values of the bubble velocity v have been given by Ref. 12. However, numerical results depend exponentially on the third power of the unknown parameter R_0, so that any answer for dN_B/dy is possible.

We can make some simple estimates of bubble number if we assume that the temperature depends only weakly on the time. Then, if the amount of supercooling before the transition occurs is small, the bubbles which are formed are deflagration bubbles and v, the bubble-growth velocity, is rather small, $v \leq 0.1$ typically.[12] In this limit, bubbles are nucleated at their critical radius but do not expand. After a nucleation time

$$\tau_N = 1/V_0(T)p(T) \ ,$$

space is filled with bubbles of a volume $V_0(T)$ and the bubble density is just $1/V_0(T)$, or

$$\frac{dN_b}{dy} = \frac{A\tau}{V_0(T)} \ . \tag{9}$$

For $A = \pi(7 \text{ fm})^2$, $\tau = 1$ fm, $V_0 =$ a few fm^3, this is a density of 10–20 bubbles per unit rapidity. The antibaryon density is some unknown fraction $\frac{1}{2}f$ of (9), before its subsequent dilution through $B\bar{B}$ annihilation.

If the supercooling is large, the bubbles are detonation bubbles, and v is the speed of sound in the medium, $v = 1/\sqrt{3}$. Then only a few bubbles are nucleated per unit volume, and the system is filled in a growth time

$$\tau_G = [3/\pi v^3 p(T)]^{1/4} \ .$$

The density of bubbles is then roughly $[1/V(T)]\tau_G/\tau_N$ for $\tau_G/\tau_N \ll 1$. This is the situation which would occur if supercooling continued until $p(T)$ were a maximum, at $\hat{T}^4 = \frac{1}{9}$. Clearly in the case of large supercooling it is easier to argue in favor of a scenario with a small bubble density per unit rapidity, while small supercooling could generate large bubble density.

Of course, the number of baryons and antibaryons which can actually escape to be detected is much less than the number which is originally made: baryons and antibaryons which are produced in the transition will annihilate against one another. Since the hadron state cools and expands rapidly after the transition, it is unlikely that all of the baryons and antibaryons can annihilate, or that the density of baryons can approach thermal equilibrium, $\propto \exp(-m_B/T)$. In particular, one expects that B's and $\bar{B}$'s produced by bubble collisions near the surface of the erstwhile plasma are more likely to survive than are those produced in the interior. In any case, as discussed by Toussaint and Wilczek,[13] conservation of baryon number during annihilation leads to production of regions of space containing mostly particles separated from regions of space containing mostly antiparticles. At long times after the phase transition, the density of particles or antiparticles is independent of the annihilation cross section and depends only on the diffusion time D and on the original density and its fluctuations.

This picture suggests several relatively model-independent consequences.

First, because baryons are topological defects, the net baryon number in a volume is equal to the topological quantum number inside the volume, which in turn is proportional to the surface integral over the volume. So the excess of baryons over antibaryons, or vice versa, is proportional to the square root of the surface area of the plasma, not its volume.[7] If after annihilation the $\bar{B}$ density is proportional to the original excess, then in the central region

$$\frac{dN_B}{dy} \propto (\text{area})^{1/2} \propto R \propto A^{1/3} \ , \tag{10}$$

where R is the nuclear radius and A is the atomic number. Alternatively, we may assume that B's and $\bar{B}$'s produced near the surface of the expanding plasma are less likely to annihilate than those on the interior. Then once again

$$\frac{dN_B}{dy} = \frac{2\pi R}{V_0} L \frac{f}{2} \tau \propto A^{1/3} \tag{11}$$

if all B's or $\bar{B}$'s produced a distance L from the surface escape, and all the ones on the interior annihilate.

Equation (11) allows a fairly unreliable estimate of the $\bar{p}/\pi^-$ ratio in a collision. Scale $dN(\pi^-)/dy$ from its CERN ISR value[14] of about 1 like $A^{2/3}$. In (11) assume $R = 1.2A^{2/3}$, $R_0 = L = \tau \approx 1$ fm. Then

$$\frac{dN(\bar{p})/dy}{dN(\pi^-)/dy} = \frac{1}{2} f \times \frac{3}{2} \times \frac{1.2}{A^{1/3}} \ . \tag{12}$$

If $f = \frac{1}{2}$, this ratio is 0.07 for uranium. Smaller f's still give a signal the same order of magnitude of conventional production mechanisms, and of course smaller A's enhance this surface effect.

Second, because it is the fluctuations in baryon number which survive the plasma, we expect to see strong correlations in phase space between BB or $\bar{B}\bar{B}$ pairs, and do not expect to observe nearby $B\bar{B}$ pairs. Azimuthal correlations occur when more B's than $\bar{B}$'s (or vice versa) were produced on one side of the hot plasma than on the other. This is markedly different from baryon production in ordinary processes, where $B\bar{B}$ pairs are correlated but BB pairs are not.

Third, there is probably an anticorrelation between energy fluctuations and $\bar{B}$ production. It may be that the transition involves relatively large supercooling followed by nucleation of a small number of detonation bubbles which explosively fill the plasma. In this case one would expect small $\bar{B}$ production but such striking signals as large transverse-energy flow $dE_\perp/dy$ with azimuthal (ϕ) symmetry (when one bubble is nucleated per unit rapidity and hits the wall of the plasma isotropically) or asymmetry (the bubble is nucleated near the surface), medium-range rapidity correlations (of particles from the same bubble) as well as possible large rapidity fluctuations caused by the collision of bubbles and subsequent local reheating of the plasma.[12] For a large number of $\bar{B}$'s we require a large number of bubbles. In that case, energy and/or angular fluctuations are likely to be small and $dE_\perp/dy\,d\phi$ flat.

On an event-by-event basis, one could imagine seeing both classes of events, as the number of bubbles fluctuates. A signal of anomalous $\bar{B}$ production *without* other events with associated rapidity fluctuations could be a signal for a second-order phase transition. Finally, the very dull scenario of uniform, nonfluctuating $dE_\perp/dy\,d\phi$ and only thermal $\bar{B}$ production is a signal that no phase transition has occurred.

This work is clearly speculative, and much remains to be done. It may be possible to give better estimates of the number of bubbles nucleated in the transition, or of the density of baryons per bubble. Nonetheless, the idea of a correlation between $\bar{B}$ production and a phase transition is sufficiently intriguing yet conservative that it is worth presenting in its present state.

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[1]For a recent review, see *Proceedings of the Fourth International Conference on Ultra Relativistic Nucleus-Nucleus Collisions, Helsinki, 1984,* edited by J. Maalampi (Springer, Berlin, in press).

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