POSSIBILITY OF MATTER-ANTIMATTER SEPARATION AT HIGH TEMPERATURE

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It is suggested, on the basis of a simplified model of nucleon-antinucleon interaction, that a spatial separation of baryons and antibaryons might occur in the blackbody radiation at very high temperatures, typically larger than $kT = 350$ MeV. Some applications to astrophysical data, like the matter density of the universe and a mechanism for quasistellar objects, are sketched.

It has been suggested by Harrison that the problems of galaxy formation and of the existence of matter in the Gamow-Lemaître cosmological model could be better understood if baryon inhomogeneities could appear in the blackbody radiation at very high temperatures.[1] In view of the large present density of matter in the universe, this effect would have to take place necessarily at temperatures kT much higher than 30 MeV. On the other hand, statistical fluctuations would never be large enough to provide this kind of effect.[2] In view of these remarks and of the fact that, at such temperatures, strong interactions are dominant, we have addressed ourselves to the following problem: Is it true that only statistical fluctuations can occur in the densities of baryons and antibaryons in a black body at high temperatures?

This problem is extremely involved. All we want to show here is that there may exist conditions which, in the framework of an oversimplified model for strongly interacting particles in thermal equilibrium, lead to a local separation between matter and antimatter.

Let us list the approximations which are made in that model:

(1) Among the strongly interacting particles which appear in the radiation, we only retain for consideration the pions and nucleons.

(2) The pions and the nucleon-antinucleon pairs constitute two independent components of the radiation. This amounts to assuming that the density matrix is diagonal in the Hilbert space of pions and pair states or, in other words, to consider the pair creations and annihilations as a small effect.

(3) The long-range interactions between nucleon and nucleon, between nucleon and antinucleon, and between antinucleon and antinucleon are described by a potential of the type known in nuclear forces which, in first approximation, can be neglected.

(4) The Gibbs potential is expanded into a virial series in the densities N/V, $\bar{N}/V$ of nucleons and antinucleons. In this expansion, only terms linear and quadratic in N and $\bar{N}$ are retained (pair-collision approximation).

(5) The short-range interaction between baryons is mainly due to the ω meson. Accordingly, it leads to a strong repulsion (hard core) between two similar baryons and to a strong attraction between baryons and antibaryons. We assume that this attraction is strong enough to bind such a pair into π, η, ρ, and ω mesons.[3] Furthermore, we assume the Levinson theorem to be valid for this nucleon-antinucleon interaction so that, in a state where there is a bound state, the phase shift decreases from π to zero when the center-of-mass momentum increases from zero to infinity.[4]

(6) We keep nonrelativistic kinematics for the motion of the center of mass of a nucleon-antinucleon pair.

(7) Temperature-dependent corrections to the interactions are neglected.

Most of these approximations are reasonable, except (2) and (4) which are very drastic. From assumption (2), the chemical potentials of nucleons and antinucleons are opposite and in fact zero in the system under consideration so that the free energy reads

$$F = F_0 + kTA(N+\bar{N}) + kTBN\bar{N}/V. \tag{1}$$

In this expression F_0 contains the pion contribution to the potential and that of free nucleons and antinucleons. N and $\bar{N}$ are the numbers of nucleons and antinucleons in a volume V.

The coefficient A is related to the pion-nucleon phase-shifts by the Beth-Uhlenbeck formula ($\hbar = c = 1$)[5]:

$$A = \sum_{I,J} \frac{(2I+1)(2J+1)}{4\pi} \int_0^\infty \ln\left(1 - \frac{\omega}{kT}\right) \frac{d\delta_I^J}{dk} dk. \tag{2}$$

Here J and I are the spin and isospin of a nucleon-pion pair, δ_I^J the corresponding phase shift, and k and ω the momentum and energy of the pion. Equation (2) is strictly valid in a static approximation for the nucleon. At large temperatures, it is easily seen that one can equivalently use Eq. (2) or add to F_0 the contribution of the lower pion-nucleon resonances treated as if they were elementary particles.[6]

The coefficient B is given by

$$B=-8\left(\frac{\pi}{mkT}\right)^{3/2}\sum_{I,J}\frac{(2I+1)(2J+1)}{16\pi}\times\int_0^\infty\frac{d\Delta_I{}^J(p)}{dp}e^{-E(p)/kT}dp. \quad (3)$$

Here m is the nucleon mass, J and I the spin and isospin of a nucleon-antinucleon pair, $\Delta_I{}^J$ the corresponding phase shift [which is real with approximation (2), otherwise it is easily shown that the summation has to be made over the eigen phase shifts]. p is the center-of-mass momentum and $E(p)$ the center-of-mass kinetic energy. Similar contributions for nucleon-nucleon pairs are not expected to give large phase shifts and are neglected.

The large and positive value of B is essential to the argument and is worth a more detailed discussion. Let us recall that one generally adds to B a term which includes the effect of nucleon-antinucleon bound states:

$$B_0=8\left(\frac{\pi}{mkT}\right)^{3/2}\sum g_n\exp(|\epsilon_n|/kT), \quad (4)$$

where ϵ_n is the binding energy. In the present case, ϵ_n is so large that the contribution of the bound states (π, η, ρ, ω) is directly written in F_0 and, according to assumption (2), is decoupled from the main effects in the nucleon-antinucleon system. Notice at this stage that there is no necessary direct connection, except via crossing, between the large binding energy and the annihilation cross section.

The B term acts like a repulsion. This can best be understood in classical terms: When a nucleon and an antinucleon have a very strong attractive interaction a large region of phase space corresponds to the making of a bound state. If the subsystem under consideration does not include bound states, there is a part of phase space which is forbidden; in this respect, nucleons look like hard spheres to antinucleons and like points to nucleons. In other words, a pair cannot exist as such at arbitrarily small distances.

In the hard-sphere picture, it is clear that nucleons and antinucleons will tend to separate spatially when the density increases, i.e., when the temperature of the black body increases. This effect is easily computed from Eq. (1) by using the chemical potentials as given by

$$0=\mu=\partial F/\partial N,$$
$$0=\bar{\mu}=\partial F/\partial\bar{N}. \quad (5)$$

The contribution to F_0 from the nucleons and antinucleons reads $kT[\varphi(N)+\varphi(\bar{N})]$ where

$$\varphi(N)=-NkT\ln[eVC(T)/N]. \quad (6)$$

In a completely nonrelativistic calculation we shall write

$$C=\sum_n g_n\left(\frac{m_nkT}{2\pi}\right)^{3/2}\exp\left(\frac{-m_n}{kT}\right)=\frac{N_0}{V}, \quad (7)$$

where the n indices indicate the resonances [in practice the nucleon and the $\Delta(33)$] with their degeneracy factor g_n. When using (7) one should leave out the A term in (1) as explained above.

It is easily seen that Eqs. (5) have only one solution for which $N=\bar{N}$ as long as BN_0/V is smaller than e. Above this value there are three solutions for N and $\bar{N}$. One is symmetric ($N=\bar{N}$), the other two solutions exchange into each other by the interchange $N\leftrightarrow\bar{N}$ and they are thermodynamically stable: They correspond to a minimum in F whereas the symmetric solution corresponds to a saddle point.

We have made numerical evaluations by writing the phase shifts as

$$\Delta_J{}^I(p)=\pi[1-4p^2/M^2] \text{ for } 4p^2<M^2,$$
$$=0 \text{ for } 4p^2>M^2, \quad (8)$$

where M is of the order of the ω mass. One finds the critical temperature to be of the order of 350 MeV. The density of nucleons is then typically of the order of the nuclear density.

Electromagnetic effects must also be taken into account. Any separation between nucleons and antinucleons tends to create regions where the density of charge is different from zero. However, electrons or positrons are attracted preferentially in these regions and make them neutral. It can be shown that the corresponding effects on the Gibbs potential are always negligible with respect to the last term in Eq. 2.

It may be instructive to consider the effect of matter-antimatter separation from a microscopic standpoint. In our model, nucleons scatter only upon antinucleons and not upon nucleons. This mechanism can keep apart a phase mostly made up of nucleons from a phase mostly made up of antinucleons. It shows how the neglect of annihilation is essential. If one reintroduces annihilation, the nucleon-antinucleon scattering cross section will be much larger than the nucleon-nucleon cross section. However, it is also found to be peaked forward so that the kinetic repulsion effect is much less obvious. It might conceivably be studied by computer techniques.

We have studied the behavior of the condensed phases in some detail. In particular, they are prevented from diffusing too rapidly because of the pressure exerted on their boundary by annihilation.

We have worked out a few astrophysical consequences of this effect if it were to take place in the Gamow universe. The main results are the following:

(1) The drops of matter and antimatter tend to grow in size by diffusion and drops of the same kind mix together. Drops of opposite type are kept apart by annihilation pressure. Taking these two effects into account, one can compute the density of matter in the universe without any free parameter and one gets the correct order of magnitude for the observed value.

(2) As is well known, one finds the cosmical radiation at 3°K.

(3) Masses of matter and antimatter which might be of the order of protogalaxies are made during the radiation era which ends at a time of one million years after the birth of the universe and a temperature of 1 eV.

(4) In view of the very early condensation of large masses of matter and antimatter, we have reinvestigated the model of a quasar as made of matter and antimatter. If, for instance, a mass of antimatter of the order of 10^{-1} to 10^{-3} the mass of a galaxy is enclosed in the center of a galaxy, it will be contracted by the annihilation pressure. Very different conditions occur if antimatter is in the form of a superstar or of a set of stars, leading however to similar observational conditions. In the second case, if one takes a density of particles of 10^7 cm^{-3} in the common boundary of matter and antimatter, a radius of 10^{17} cm, and a temperature of 30 000°K, one easily computes the diffusion at the boundary. It is found that the amount of energy released is 10^{47} ergs/sec; the region of annihilation has a width of 10^{12} cm and is presumably quite unstable. The energy is released mainly into photons, electrons, and positrons. Within the center, these particles are absorbed and heat antimatter. The mean free path of outgoing photons is of the order of 10^{20} cm and photons give rise mainly to electrons (by Compton effect on matter in the outskirts). These electrons are produced mainly in the forward direction so that coherent lines of current are produced which can easily generate very high magnetic fields. The radio production takes place mainly by synchrotron radiation from the electrons and positrons of annihilation in a large region with a size of the order of that of a galaxy. The radio spectrum is therefore what has been computed by Ekspong, Yamdagni, and Bonnevier.[7] The lifetime of a quasistellar object depends on the mass of its antimatter center but can easily attain several times 10^{10} yr. Obviously the three main objections to this kind of models: "Why antimatter, why not hard photon emission, why so high magnetic fields?" are naturally met with in this version.

More details will be published elsewhere.

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[3]This idea was first suggested by Fermi and Yang for pions. For a recent treatment see J. S. Ball, A. Scotti, and D. Y. Wong, Phys. Rev. 142, 1000 (1966).

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