Cosmological model of a baryon island

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The observational data seem to indicate that the visible matter in our universe is concentrated inside a sphere with radius corresponding to red shift $z \approx 5$. A model which explains this phenomenon is presented. The cosmological evolution of the baryon island and the observational consequences of this hypothesis are discussed.

1. INTRODUCTION

In cosmology the dependence of any actually observable large-scale property of the universe on the red shift z is usually interpreted either as a selection effect of the observations or as an evolutionary effect associated with the direction of the development of the whole universe. The best-known example is the absence of quasars at large z (Ref. 1). However, recent astronomical data appear to suggest the genuine absence of visible matter at large spatial distances exceeding $z \approx 5$.

Thus, despite careful searches[2,3] for small-scale anisotropy $\Delta T/T$ in the background relic radiation, this anisotropy has not yet been detected. This is hard to explain in the framework of the standard cosmological model,[4] in which the anisotropy $\Delta T/T$ is closely connected with matter fluctuations at large z. As we shall see below, the absence of small-scale anisotropy can be explained in the framework of the assumption that ordinary matter is absent at large distances. Of course, the interpretation of these data is ambiguous, and more-sensitive observations are necessary. However, if the above tendency continues in the near future, and with a higher level of accuracy of the measurements, it will be possible to assert that at large distances ordinary matter in the form of protons, nuclei of heavier elements, and electrons has an appreciably smaller density than in the part of the universe nearer to us. In other words, it can be conjectured that our universe is a baryon island[5] in a sea of some kind of invisible matter (or, perhaps, even in a vacuum).

Such baryon (and antibaryon) islands can arise in a theory with spontaneous violation of *C*- (and *CP*-) invariance, a necessary condition being that the process of spontaneous violation occur during the inflationary stage not long before the end of this stage.[6] The latter condition ensures that the island is sufficiently large but, nevertheless, has its "shore" inside the present-day horizon.

As a result of processes that we shall consider in detail in Sec. 2, in the stage under consideration regions (bubbles) arise inside which there is a small excess of baryons over antibaryons (or vice versa) and outside which the universe is completely charge-symmetric. In the expansion and cooling of the universe the baryons and antibaryons in the symmetric regions annihilate each other almost completely; the residue turns out to be negligibly small.

According to standard estimates in the framework of a charge-symmetric model, the relation $N_b = N_{\bar{b}} \approx 3 \cdot 10^{-19} N_\gamma$ is valid for the densities of baryons, antibaryons, and electromagnetic radiation.[7] In the proposed scheme the excess of baryons (and, correspondingly, of electrons) inside a bubble does not vanish and is equal to $N_b \approx 4 \cdot 10^{-10} N_\gamma$. As a result, separate baryon (or antibaryon) islands arise on a background of charge-symmetric matter. The distance between the individual islands, and their sizes, depend strongly on the parameters of the model. In a favorable case we may expect that our island has its shore within the visibility horizon; the existence of other islands within the visibility horizon is also possible.

If the island model under consideration is valid, the relic radiation reaching us today originates from regions outside the island, where baryons are absent, and, therefore, fluctuations of the temperature of the radiation are not connected with variations in the observable density of baryons, as is the case in the standard theory of the formation of the large-scale structure of the universe. Because of the absence of baryons in the space between the islands, the relic radiation there is decoupled from matter not at $z = 10^3$, as in the usual homogeneous model, but at $z = 10^7$, when the optical depth with respect to Thompson scattering of the relic radiation by the "quenched" electron-positron pairs that have remained after the annihilation[1)] ($N_{e^-} \approx N_{e^+} \approx 10^{-6} N_\gamma$) gradually becomes small ($\tau_T < 1$).

An asymmetric (noncentral) position of the observer in our island could lead to large-scale asymmetry of the relic radiation. The observed small asymmetry of the relic radiation imposes a substantial restriction on the position of the observer in the island. The main possibility of experimental confirmation involves making observations of objects near the boundary of the island, and, possibly, detecting other, more distant islands (although the latter is unlikely).

The arrangement of the material of the article is as follows. In Sec. 2 we construct a model of spontaneous *C*- (*CP*-) violation with the properties we need. In Secs. 3 and 4 we discuss possible astronomical observations for a test of the island model.

2. MODEL OF SPONTANEOUS *C*- (*CP*-) VIOLATION

A. General scenario

In its basic features the scenario under consideration is as follows. The *C*- (and *CP*-) violation is connected with the complex condensate of a scalar field φ. The transition from the symmetric phase with expectation value $\langle \varphi \rangle = 0$ to a *C*- and *CP*-odd phase with $\langle \varphi \rangle \neq 0$ (with $\mathrm{Im} \langle \varphi \rangle \neq 0$) occurs dynamically as a consequence of the interaction of the field φ with the inflaton field Φ responsible for the inflation, when the condensate $\langle \Phi(t) \rangle$ slowly slides to its equilibrium value Φ_0. We shall assume that this interaction has the form $\lambda_\Phi |\varphi|^2 (\Phi - \Phi_1)^2$, where $\Phi_1 < \Phi_0$. Obviously, the transition

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to the CP-odd phase is more probable when $\Phi \sim \Phi_1$.

The scenario under consideration can also be realized with another form of interaction between φ and Φ, though perhaps not so naturally. In the process of the evolution, when Φ is close to Φ_1, in the universe separate regions with $\langle \varphi \rangle \neq 0$, in which the C- (CP-) symmetry is spontaneously broken, are formed on the background of the CP-even part of the universe with $\langle \varphi \rangle = 0$. At the time of formation the size of these C- (CP-) odd bubbles is equal to $d_0 \propto m_\varphi^{-1}$, where m_φ is the mass of the field φ. In order that our observable universe, with a size corresponding to $z \approx 5$, grow from such a bubble, a period of exponential expansion (inflation) lasting approximately $70H^{-1}$ is necessary, where H is the Hubble constant at that time. After the end of the inflation secondary heating occurs and a high-temperature plasma of elementary particles is formed. The state of the matter outside and inside the bubble at this time is approximately the same, the sole difference being that outside the bubble the interactions of particles and antiparticles are fully symmetric, while inside the bubble there is a certain (small) asymmetry. The further evolution of the model will be considered in Sec. 3.

B. Spontaneous violation of charge invariance; basic features of the model

A model of spontaneous violation of C- (CP-) invariance was first proposed in Ref. 8. As is usual in spontaneous symmetry breaking, the Lagrangian is symmetric while the ground state (vacuum) is not. The vacuum in such a theory is degenerate, i.e., there exist several vacuum states, related to each other by symmetry transformations. In the simplest model, however, the phase transition from the symmetric to the asymmetric phase that occurs as the temperature decreases[9] is a second-order phase transition, occurring without supercooling and without the formation of bubbles of the new phase in the old phase. Here the asymmetric phase is formed practically simultaneously over the entire space and a domain structure arises, in accordance with the sign of the amplitude of the C- (CP-) violation ($\mathrm{Im}\langle \varphi \rangle > 0$ or $\mathrm{Im}\langle \varphi \rangle < 0$). In summary, regions with an excess of baryons or antibaryons, bordering on each other and separated by domain walls with high energy density, arise in the universe.

In order to obtain individual baryon or antibaryon islands that are sufficiently remote from each other it is necessary to construct a theory in which the state with $\langle \varphi \rangle = 0$ is metastable and the transition to the C- (CP-) odd vacuum with $\langle \varphi \rangle \neq 0$ is a first-order phase transition. This can be achieved if the potential of the field φ has the form

$$v(\varphi) = m_1^2 |\varphi|^2 + \frac{1}{2} m_2^2 (\varphi^2 + \varphi^{*2}) + \frac{\lambda}{2} |\varphi|^4 \ln\left(\frac{|\varphi|^2}{\eta^2}\right), \quad (2.1)$$

where η is a constant of the problem, the masses characterizing the field satisfy the inequality $m_1^2 > m_2^2 > 0$, the constant λ is positive (see Fig. 1), and φ^* is the complex conjugate of the field. The logarithmic dependence on φ, as is well known, arises when one-loop radiative corrections to the self-interaction $\lambda |\varphi|^4$ are taken into account.[10]

The probability of tunneling of the field φ from the state $\varphi = 0$ to a deeper minimum of the potential ($|\varphi_0|^2 = \eta^2/e^{1/2}$ for $m^2 \ll \lambda \eta^2$) in the quasiclassical approximation is determined by the factor $\exp(-C/\lambda)$, where $C \gtrsim 1$. To be precise, for the model under consideration we have

$$C = 8\pi^2/\{3 \ln [\lambda \eta^2 (m_1^2 - m_2^2)^{-1}]\}. \quad (2.2)$$

Since, typically, $\lambda \ll 1$, the probability of a phase transition to a C- (CP-) odd phase is negligible. However, if the field φ interacts with the field Φ (the inflaton field) responsible for the inflation, it is not difficult to construct a model in which the probability of spontaneous violation of C (CP) is large, specifically in the final stage of the inflation.

C. The inflaton—the generator of nuclei of the *CP*-odd phase

We shall assume that the inflation scenario of Ref. 11 obtains, according to which exponential expansion occurs when, and on account of the fact that, Φ is far from its equilibrium value Φ_0, the rate of change of $\Phi(t)$ being small in comparison with the rate H of expansion of the universe. The form of the potential energy $V(\Phi)$ is depicted schematically in Fig. 2. As Φ approaches Φ_0 the potential becomes steeper, Φ begins to change more rapidly, and the law of the expansion changes from an exponential to a power law. Oscillations of Φ about the equilibrium value Φ_0 lead to the creation of particles and to (secondary) heating of the universe.[12] In the standard scenario the baryon asymmetry of the universe arises at this time. The time dependence of $\Phi(t)$ is depicted qualitatively in Fig. 3.

Let the interaction between the fields φ and Φ have the form

$$L_{int} = \lambda_\Phi |\varphi|^2 (\Phi - \Phi_1)^2, \quad (2.3)$$

where $\Phi_1 < \Phi_0$ and $\lambda_\Phi > 0$. This interaction gives rise to an effective, time-dependent mass of the field φ:

$$m_\varphi^2(t) \approx \lambda_\Phi [\Phi(t) - \Phi_1]^2.$$

We shall assume that the parameters are chosen in such a way that $m_1^2 + m_\varphi^2(t) > m_2^2 > 0$ when $(1 - \Phi/\Phi_1) \gtrsim O(1)$, but $m_2^2 > m_1^2$. Then the state $\varphi = 0$ will be classically stable against small disturbances almost all the time, and only for $\Phi \approx \Phi_1$ will there be a period of instability. Quantum fluctuations of the field φ grow at this time, and if they are able to exceed a certain critical value φ_c (see Fig. 1) by the time

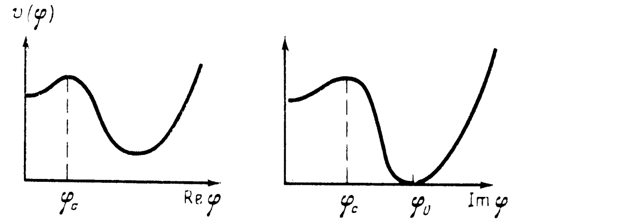

FIG. 1. Effective potential of the field φ in one-dimensional projection on the plane (v, Reφ) and (v, Imφ). The case when the potential reaches its minimum value at Re$\varphi = 0$ and Im$\varphi \neq 0$ is depicted.

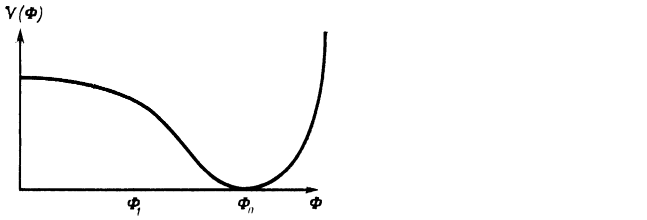

FIG. 2. Effective potential of the inflaton field; Φ_0 is the value of Φ at which $V(\Phi)$ is a minimum.

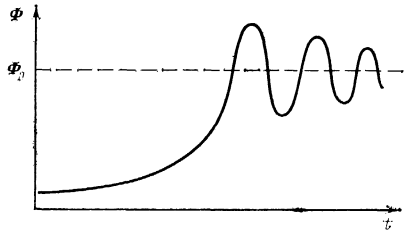

FIG. 3. Time dependence of the inflaton field.

when the condition $m_1^2 + m_\varphi^2(t) > m_2^2$ is again fulfilled and the condition $\varphi = 0$ again becomes stable, the disturbances that have grown will no longer return to $\varphi = 0$ but will grow to $\varphi = \pm i\varphi_0$. As a result, superposed on the charge-symmetric world there arise spatial regions in which C- (CP-) invariance is violated. The average size d of the C- (CP-) odd bubbles and the distance l between them depends strongly on the parameters of the model, and, in particular, on the duration of the period of instability.

The equation of motion of the inflaton is modified on account of the interaction (2.3):

$$\ddot{\Phi}+3H\dot{\Phi}+U_\Phi'(\Phi)+\lambda_\Phi|\varphi|^2(\Phi-\Phi_1)=0\ . \qquad (2.4)$$

Here a dot denotes a time derivative and a prime denotes a derivative with respect to the field Φ. This modification will be of little importance if $\lambda_\Phi|\varphi|^2$ is small in comparison with U'_Φ. However, the analogous term can have an important effect on the evolution of the field φ. The equation of motion satisfied by the values of φ inside the C-odd bubble has the form

$$\ddot{\varphi}-\frac{1}{a^2}\Delta\varphi+3H\dot{\varphi}+\frac{dv}{d\varphi^*}+\lambda_\Phi[\Phi(t)-\Phi_1]^2=0, \qquad (2.5)$$

where a is a scale factor and $\Delta\varphi$ is the Laplacian of the field φ. It is not difficult to analyze the solutions of this equation qualitatively for the long-wavelength modes of φ, when we can neglect the term containing $\Delta\varphi$. The result is that φ tends to an equilibrium value with possible oscillations about the equilibrium position. We note the possibility of a parametric resonance associated with oscillations of $\Phi(t)$.

D. Density inhomogeneities in an island universe

An interesting feature of the model under consideration is the fact that variations in the distribution of the visible (baryonic) matter are not connected with temperature fluctuations of the relic radiation, which presently reaches us from a region in which baryons are absent. This makes it possible to have rather large baryon-density fluctuations, unrestricted by the magnitude of $\Delta T/T$, at the time of recombination. Moreover, for the generation of significant inhomogeneities $\delta\rho_b$ there exists in the model a natural mechanism associated with spatial fluctuations of the amplitude of the C- (CP-) violation. As we know, the latter amplitude is proportional to the imaginary part of the condensate $\langle\varphi(x,t)\rangle$. It is obvious that in this case isothermal density perturbations arise, in which N_b/N_γ varies from point to point while N_γ remains constant (N is the particle-number density).

It is well known[13] that quantum fluctuations of the scalar field grow strongly as a consequence of inflation. There then arises an almost flat (to within logarithmic terms) fluctuation spectrum with amplitude $\sim 100\lambda^{1/2}$ (for the $\lambda\varphi^4$ theory). Usually it is assumed that these fluctuations of the scalar field generate the observable density inhomogeneities; therefore, in order to obtain $\delta\rho/\rho \sim 10^{-4}$ it is assumed that λ is unnaturally small: $\lambda \approx 10^{-12}$. In our case fluctuations of φ lead to fluctuations in the magnitude of the baryon asymmetry but not in the total density of matter, which consists principally of nonbaryonic matter. Hence there is no need for such a small value of λ.

In the model under consideration there is one further mechanism for generating density perturbations with a spectrum possessing a sharply pronounced maximum in the region of large wavelengths. In order to clarify how such inhomogeneities arise, we shall consider the following simplified model. Suppose that the initial fluctuation of the field φ, from which a CP-odd vacuum condensate $\pm i\varphi_0$ later develops, has the form depicted in Fig. 4. We recall that φ_c and L_c are the critical values of the field φ and fluctuation wavelength at which the state near $\varphi = 0$ becomes unstable. The potential energy of the field φ, when φ is far from the equilibrium position φ_0, can be written approximately in the form

$$v(\varphi) \approx \frac{1}{2}m_1^2|\varphi|^2+\frac{\lambda}{2}|\varphi_1|^4\ln\frac{|\varphi^2|}{\eta^2}. \qquad (2.6)$$

Here we consider "sliding" along the purely imaginary direction in the complex plane of φ, and, therefore, we can regard φ in this expression as real. It is not difficult to see that the quantity φ_c is connected with the parameters of the potential (2.6) by the relation

$$\varphi_c^2=\frac{m_1^2}{2\lambda\ln(2\lambda\eta^2/m_1^2)} \qquad (2.7)$$

(for $\lambda\eta^2 \gg m_1^2$).

During the inflation period the motion of φ toward the equilibrium position can be assumed to be slow because of the large force of "Hubble friction" (the term $3H\dot{\varphi}$ in the classical equation of motion). For a qualitative analysis of the phenomenon we shall neglect the expansion of the universe and assume that φ satisfies the equation of motion

$$\ddot{\varphi}+v_\varphi'=0 \qquad (2.8)$$

with the initial condition $\varphi(t=0) = \varphi_c f(r)$, where $f(r)$ is the function describing the shape of the initial fluctuation of φ, depicted in Fig. 4. Here r is regarded as a parameter, and $f(r)$ is not necessarily spherically symmetric. For simplicity we shall assume that $\dot{\varphi}(t=0) = 0$. As long as the value of φ is still far from the equilibrium position φ_0 but already far from φ_c as well, the solution of Eq. (2.8) has the form

$$\varphi=\lambda^{-1/2}[t_0(r)-t]^{-1}. \qquad (2.9)$$

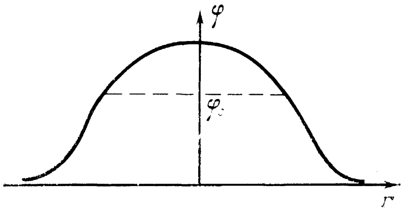

FIG. 4. Schematic shape of an initial fluctuation of the field φ. The length of the dashed line is equal to L_c.

For φ close to φ_c the solution has an oscillatory character:

$$\varphi=\varphi_0\{1+A\sin m_2[t-t_0(r)-\lambda^{1/2}\varphi_1^{-1}]\}, \qquad (2.10)$$

where φ_1 is the value of φ corresponding to the crossover from the regime (2.9) to the regime (2.10), and A is the amplitude of the oscillations.

The quantity $t_0(r)$ can be related to the initial field distribution $\varphi(t=0, r)$. Using for this the formula (2.9), although this is not completely rigorous, we find

$$t_0(r)=\lambda^{1/2}\varphi_c^{-1}/f_0(r). \qquad (2.11)$$

Now, by means of the expression (2.7), it is not difficult to show that the characteristic size of the spatial fluctuations of the amplitude of the C-(CP-) violation is equal to

$$l_c\simeq\left(\frac{m_1}{m_2}\right)^2 l_u/\delta f(r), \qquad (2.12)$$

where l_u is the present-day size of our baryon island and δf is the variation of the function $f(r)$. In the usual models $m_1 \ll m_2$, and, e.g., for $m_1 = 10^{-3} m_2$ and $\delta f \cong 0.1$ we obtain $l_c = 10^{-2} l_u \approx 10^8$ yr.

In a realistic situation, if the parameter m_2 is large in comparison with the Hubble constant at the time of formation of the baryon asymmetry, the spatial and temporal oscillations of φ will vanish upon time averaging. However, if at this time $m_2 \approx H$, then a periodic (or almost periodic) distribution of baryons, with characteristic size l_c (2.12), can arise in the universe.

E. Evolution of the boundary of the island

In accordance with the proposed model, the universe as a whole consists of separate baryon or antibaryon islands—mini-universes, immersed in invisible C- (CP-) symmetric matter. During the period of formation of the baryons the size was determined by the size of the region in which C-(CP-) invariance was spontaneously violated, i.e., a condensate $\langle\varphi\rangle \neq 0$ arose. Outside, there is no field condensate and the state φ is metastable. The difference in the energy densities of the true ($\langle\varphi\rangle = \pm i\varphi_0$) and spurious ($\langle\varphi\rangle = 0$) vacua is equal in order of magnitude to $\rho_{\varphi_0} = \lambda\varphi_0^4$, and the thickness of the transitional region is $\propto m^{-1}$ (where m is a characteristic mass parameter). By assumption, $\rho_{\varphi_0} \ll \rho_{\text{tot}}$ in the period of baryosynthesis.

After the formation of baryons ceases with decrease of the temperature, the boundary of the baryon island expands, together with the universe, in accordance with the Friedmann law $a \propto t^{1/2}$ or $a \propto t^{2/3}$, while the walls of the bubble of the true vacuum move in the spurious vacuum with almost the speed of light (if we neglect the possible interaction of the field φ with the primary plasma). Therefore, the size of the C- (CP-) odd bubble becomes greater than the size of the baryon island.

By the time ρ_{tot} is falling sufficiently strongly, so that ρ_{tot} becomes of order ρ_{φ_0}, the law of expansion of the universe outside a C- (CP-) odd bubble should become exponential, since, by assumption (or, more precisely, according to astronomical observations), the vacuum energy inside the bubble is equal to zero while that outside the bubble is equal to ρ_{φ_0}. As a result, the energy density of all forms of matter except φ outside the bubble practically vanishes. Therefore, the relic radiation in our mini-universe should also vanish.

There are several ways of tackling this contradiction. First, it is possible to suppose that the same (at present, unknown) mechanism that annihilates the cosmological constant in our world also annihilates it beyond the limits of our world. Then the laws of expansion outside and inside the bubble should not differ appreciably.

The second possibility is that the region with $\langle\varphi\rangle = 0$ expands exponentially with $\rho_{\text{vac}} = \lambda\varphi_0^4$, and the relic radiation reaches us from the walls of the bubble which process the energy of the spurious vacuum into energy of particles. In this case the plasma that is formed should have time to thermalize, in order that the arriving radiation have a Planck spectrum.

Finally, the third possibility is that the difference of the energies of the field φ inside and outside the bubble is dissipated as a result of a mechanism of the type considered in Ref. 14. This mechanism can be realized if in the effective mass of the field φ we take into account an additional term of the form

$$\delta m^2=aT^2+bR, \qquad (2.13)$$

where T is the temperature and R is the curvature scalar. Here the square of the mass of φ for small values of the field can be negative (when T and R become sufficiently small), and a second-order phase transition to a true vacuum occurs throughout space.

It is possible to consider another variant, in which spontaneous CP-violation remains even after a second-order phase transition to $\langle\varphi\rangle = \pm i\varphi_0$. Then the region outside the large C- (CP-) odd bubble in which our mini-universe is situated will have a cellular structure, in which regions with $\varphi = \pm i\varphi_0$ are mixed with regions with $\varphi = -i\varphi_0$ and are separated by domain walls. The characteristic size of such regions will not be large, and therefore their presence will not lead to appreciable anisotropy. However, in this case the energy contained in the walls turns out to be too large,[15] if we do not assume an unnaturally small value for the characteristic mass scale of the field φ, i.e., $m_\varphi \leqslant 0.1$ MeV.

In the light of the above, the most natural model seems to be one with dynamical cancellation of the Λ-term, although a concrete example of such a model has not yet been found. In its favor, however, is the fantastic accuracy of the cancellation of the Λ-term in comparison with the theoretical expectation. The construction of such a model will make it possible to solve at the same time the domain-wall problem that faces us.

In any case, for this picture of the formation of an island universe to be viable one of the above mechanisms (or some other mechanism) of annihilation of domain walls must operate.

3. DYNAMICAL EVOLUTION OF THE MODEL

We shall consider the dynamical evolution of the matter of the universe outside and inside a baryon island. We set the quantity $\Omega \equiv \rho/\rho_{\text{cr}}$ equal to unity. For simplicity we shall assume that the island has the shape of a sphere and that the density of the excess of baryons is uniform everywhere within the island. The dimensions of the island are such that today they are comparable to the distance to the horizon. In the past the dimensions of the island were much greater than the distance to the horizon, and initially, in the first approximation, we shall neglect all boundary effects, assuming that the matter inside and outside evolves independently. We shall consider boundary effects later.

The dynamics of the expansion of matter outside and inside the island becomes different after the annihilation of baryon-antibaryon pairs has occurred and the excess of baryons remaining inside the island has become nonrelativistic. We denote this time by t_0. In order of magnitude, $t_0 \approx 10^{-6}$ sec, and the temperature T_0 at this time corresponds to the baryon mass m_b. The total density of all forms of matter at this time, both inside and outside the island, is equal to[7]

$$\rho_0 = 3/32\pi G t_0^2, \tag{3.1}$$

where G is the gravitational constant. After this time we divide all matter into three categories: baryons with mass m_b (together with an equal number of electrons), invisible matter consisting of light weakly interacting particles (conventionally called neutrinos) with $m_\nu \ll m_b$, and photons. For the estimates we set $m_\nu \approx 10$ eV. At the time t_0 we denote the distribution of the density between the components as follows:

$$\rho_b = \alpha\rho_0, \quad \rho_\nu = A\rho_0, \quad \rho_\gamma = B\rho_0, \quad \alpha + A + B = 1. \tag{3.2}$$

For order-of-magnitude estimates we can set $A \approx B \approx \frac{1}{2}$ and $\alpha \approx N_b/2N_\nu \approx 2\cdot 10^{-10}$ inside the island and equal to zero outside it; here $N_\nu \approx A\rho_0/m_\nu$. The neutrino component remains relativistic up to the time t_1 when the temperature T_1 corresponds to m_ν. In order of magnitude, $t_1 \approx (m_b/m_\nu)^2 t_0 = 10^{10}$ sec [for a more accurate expression, see (3.4)].

The solution of the equations of gravitation for the mixture of baryons without pressure ($p = 0$) and the relativistic component ($p = \rho/3$) can be written in the form

$$a = \left(\frac{t}{t_0}\right)^{1/2}\left[1 + \frac{\alpha}{6}\left(\frac{t}{t_0}\right)^{1/2}\right], \tag{3.3}$$

where the time t is the proper time of a comoving observer and a is the scale factor. It is assumed that $(\alpha/6)(t/t_0)^{1/2} \ll 1$. For our model the solution is applicable in the interval $t_0 < t < t_1$. The solution (3.3) is applicable both inside the island, where $\alpha \neq 0$, and outside it, where $\alpha = 0$. The rates of change of a outside and inside the island are slightly different, and, therefore, the same temperature T_1, corresponding to m_ν, is reached at slightly different times

$$t_1 = t_0 (m_b/m_\nu)^2 \left[1 - \frac{\alpha}{3}\left(\frac{m_b}{m_\nu}\right)\right]. \tag{3.4}$$

After the time t_1 (which is different outside and inside the island!) the neutrino component becomes nonrelativistic ($p_\nu = 0$). In the subsequent evolution the gravitation of the photons can be neglected. The solution of the equations of gravitation for $t > t_1$ can be written in the form

$$a = [6\pi G A\rho_0 (m_\nu/m_b)]^{1/3}\left(1 + \frac{\alpha m_b}{3A m_\nu}\right) \cdot \left\{t - t_0\left(\frac{m_b}{m_\nu}\right)^2\left[1 - \frac{4}{3A^{1/2}} - \frac{\alpha}{3}\frac{m_b}{m_\nu}(1 - 2A^{-1/2})\right]\right\}^{2/3}. \tag{3.5}$$

As before, inside the island $\alpha = 2\cdot 10^{-10}$, and outside the island $\alpha = 0$. The nonrelativistic density component $\rho_\bullet$ inside the island consists of baryons and neutrinos, while that outside the island consists only of neutrinos. For $t \gg t_1$ we have for the difference in the densities of the nonrelativistic components inside and outside the island and for the corresponding differences of the Hubble constant (as before, we assume that the boundary zone of the island, where the motion is perturbed, is small)

$$\frac{\Delta\rho_\bullet}{\langle\rho_\bullet\rangle} \approx \frac{2}{3}\left(\frac{m_b}{m_\nu}\right)^3 \alpha \frac{t_0}{t}(2A^{-1/2} - 1), \tag{3.6}$$

$$\frac{\Delta H}{\langle H\rangle} = \frac{1}{2}\frac{\Delta\rho_\bullet}{\langle\rho_\bullet\rangle}. \tag{3.7}$$

For the difference of the energy densities of the photons inside and outside the island, up to times when the size of the island is still much greater than the distance to the horizon but $t \gg t_1$, we have

$$\frac{\Delta\rho_\gamma}{\langle\rho_\gamma\rangle} = -\frac{4}{3}\frac{\alpha}{A}\frac{m_b}{m_\nu} \approx -\frac{4}{3}\Omega_b \approx -0.04, \tag{3.8}$$

where Ω_b is the ratio of the present-day baryon-energy density to the critical baryon-energy density.

We shall make a few remarks concerning effects at the edge of the island. During the period $t_0 < t < t_1$ the edge of the island expands more rapidly than the surrounding matter; see (3.3). The relative velocity of the edge and the surrounding matter is of the order of the speed of light. This leads to the result that an excess of mass is formed at the edge of the island. In this case neutrinos pass freely through the matter, and hydrodynamic forces act on the photons and baryons. The excess mass at the edge is built up principally at $t \approx t_1$, when a change in the equation of state occurs. With neglect of nonlinear effects, the relative mass excess $\Delta M/M$ at the edge of the island is equal to the ratio ρ_b/ρ_ν of the baryon density to the density of hidden mass (an estimate is easily obtained by considering the properties of the expansion, as $t \to \infty$, of a sphere with the critical density and $p = 0$). This conclusion is valid not only in the light-neutrino model under consideration, but also when the particles of hidden mass have a rest mass much greater than the rest mass of the baryons.

It is probable that boundary effects reduce the value of $\Delta M/M$ somewhat, and therefore, in order of magnitude, we take

$$\Delta M/M \approx 10^{-2}. \tag{3.9}$$

In the narrow boundary spherical layer we have $\Delta\rho/\rho \approx 1$, and the behavior of the matter requires separate study. It is clear that the conditions of the formation of the structure here will be different from those inside (and outside) the island. We note that the picture we have painted of the structure of the boundary region depends on our assumption of a sharp boundary to the averaged baryon-density distribution in the island. For a smooth distribution of baryons the picture will be different. If in the initial conditions, at $t = t_0$, the density is everywhere equal to the critical density, the excess of mass in the baryon island means that in the future the expansion will give way to compression.

4. OBSERVATIONAL VERIFICATION

Evidently, an observational test should involve exhibiting the boundary of our island, determining our position relative to the center of the island, investigating the matter beyond the boundary, and attempting to detect other islands.

A. Observations of the anisotropy of the relic radiation

We shall consider systematically various effects that lead to anisotropy of the relic radiation. We shall assume that the position of the observer is not at the center of the island. Let the distance to the farthest edge correspond to red shift z_* and the distance to the nearest edge correspond to red shift z_{**}. We shall begin with effects that are not connected with the presence of an excess of mass on the boundary of the island. The rates of expansion inside and outside the island are slightly different (see the preceding section). Therefore, when a ray of relic radiation coming toward the observer from the horizon that lies in the direction of the farthest edge of the island has already entered the island, while a ray coming from the opposite direction is still outside it, their reddening will be somewhat different. This will create large-scale anisotropy. From formula (3.5) it is possible to obtain an expression for the difference in the temperatures of the relic radiation in the directions of the farthest and nearest edges of the island (t_c is the present time):

$$\Delta T/T = -\frac{2}{9}\frac{t_0}{t_c}\alpha\left(\frac{m_b}{m_\nu}\right)^3 (2A^{-1/3}-1)\left[(z_*+1)^{3/2}-(z_{**}+1)^{3/2}\right]. \tag{4.1}$$

This formula is applicable for $t_* \gg t_1 \simeq 10^{10}$ sec (i.e., $z_* \ll 10^5$).

If for the estimate we take $z_* = 8$ and $z_{**} = 3$, then $\Delta T/T = 3\cdot 10^{-8}$.

In view of the nonstationary character of the island, a much greater anisotropy $\Delta T/T$ is induced by the presence of the gravitational effect of the mass excess ΔM on the edge of the island in the model under consideration.

For values of z_* and z_{**} of the order of a few units, when the Newtonian-potential approximation is applicable, we obtain the following formula:

$$\frac{\Delta T}{T} = \frac{1}{2}\frac{\Delta M}{M}(z_*-z_{**})\left[2-\frac{1}{(z_*+1)^{1/2}}-\frac{1}{(z_{**}+1)^{1/2}}\right]^2$$

$$\approx (z_*-z_{**})/200. \tag{4.2}$$

It should be emphasized that the principal harmonic of $\Delta T/T$ is the dipolar harmonic. Therefore, the restriction on $z_* - z_{**}$ that follow from the observations is $z_* - z_{**} < 0.2$. We note also that the estimates given for $\Delta T/T$ are model-sensitive to the degree of nonsphericity of the island and to the distribution of the averaged baryon density inside the island.

For the anisotropy of the relic radiation we shall make a further estimate, connected with the possibility of scattering of the radiation by the intergalactic plasma. Suppose that inside the baryon island there is a plasma constituting a fraction $\Omega_{\rm pl}$ of the total density of matter. Then the optical depth with respect to scattering out to an island boundary corresponding to red shift z is equal to (see, e.g., Ref. 7)

$$\tau = 0.03\left(\frac{H_0}{75}\right)\Omega_{\rm pl}\left[(1+z)^{3/2}-1\right], \tag{4.3}$$

where H_0 is the present-day value of the Hubble constant in km sec^{-1} Mpc^{-1}. For $\tau \ll 1$, because of scattering by plasma mainly at the edge of the island (where the rays were to be found in the past, at high density) the intensity of the rays coming directly toward the observer from beyond the limits of the island is weakened by a factor of $1-\tau$.

On the other hand, the scattering leads to the result that there begins to propagate in the direction toward the observer a fraction τ of the rays that traveled originally, on the boundary of the island, in other directions. Of these, about half come from the region inside the island, and have lower temperature, while the other half come from regions outside the island. As a result, a crude estimate of the radiation temperature, as measured from the Rayleigh-Jeans region of the spectrum, is given by the expression

$$T = (1-\tau)T_1 + \frac{\tau}{2}T_2 + \frac{\tau}{2}T_1 = \left(1-\frac{\tau}{2}\right)T_1 + \frac{\tau}{2}T_2, \tag{4.4}$$

where T_1 and T_2 are the temperatures of the photons outside and inside the island, respectively. Hence, for the anisotropy associated with the scattering we obtain

$$\Delta T_{\rm sc}/T \simeq \frac{\tau_*-\tau_{**}}{2}(\Delta T/T)_{max}. \tag{4.5}$$

Here τ_* and τ_{**} are the optical depths in the directions of the furthest and nearest edge of the island, respectively, and $(\Delta T/T)_{\rm max} = 1/4\Delta\rho_\gamma/\rho_\gamma = 0.01$ [see (3.8)]. Substituting the expression (4.3) into (4.5), we obtain (for $H_0 = 75$ km sec^{-1} Mpc^{-1}),

$$\Delta T_{\rm sc}/T \approx 10^{-4}\Omega_{\rm pl}\left[(z_*+1)^{3/2}-(z_{**}+1)^{3/2}\right]. \tag{4.6}$$

In the case $z_* > z_{**}$, in the direction of the furthest boundary of the island a broad (of angular size $\omega = 2\pi$ sr) spot with a slightly lowered temperature will be observed. The spherical harmonics (dipolar, quadrupolar, etc.) of such a spot should be correlated in direction and have comparable amplitudes. Their absence (apart from terms of order $\Delta T/T < 10^{-4}$) leads (for $z_* \gg z_{**}$) to a bound on the quantity of plasma in the island:

$$\Omega_{\rm pl} < z_*^{-3/2} \approx 0.1\,(z_*/4)^{-3/2}. \tag{4.7}$$

We turn to the small-scale relic-radiation anisotropy $\Delta T/T$ associated with the presence of a density perturbation $\delta\rho/\rho$. In our model, as already noted, there are at least two possible ways in which large-scale structure inside the island can be formed.

The first possibility presupposes a scenario analogous to the standard model of a homogeneous universe. In this case, in the early stages, there were adiabatic perturbations $\delta\rho/\rho$ of the matter everywhere—both inside and outside the island. In this case, $\delta\rho/\rho$ outside the island will lead to fluctuations $\Delta T/T$ on account of the Sachs/Wolfe effect.[16] The Silk effect[17] and Syunyanev–Zel'dovich effect[18] are absent here over all the angular scales accessible to observation, since the universe is transparent out to $z \approx 10^7$ instead of the value $z \approx 10^3$ in the standard model. The large transparency leads to a logarithmic growth of $\Delta T/T$ over small angular scales (tens of minutes) on account of the Sachs–Wolfe effect, but the absence of other effects (which have the same order of magnitude in the standard theory) leads to the result that, as a whole, this anisotropy will remain of the same order of magnitude as in the standard theory.

According to the second possibility of structure formation, we can assume that the fluctuations $\delta\rho/\rho$ outside the island are arbitrarily small and the structure inside the island is formed as a result of large spatial variations of baryon charge. In this case the small-scale anisotropy of the fluctuations $\Delta T/T$ of the relic radiation can turn out to be very small, since the anisotropy comes from outside the island, where there are practically no appreciable fluctuations $\delta\rho/\rho$.

B. Background annihilation radiation

The appearance of additional background radiation in the island model is connected with processes following the annihilation of electron-positron pairs. In fact, the universe beyond the boundary of the baryon island, at $z > z_{\rm bound} \approx 4$, is not entirely free of baryons and charged leptons. In the case of a charge-symmetric universe, its volume is occupied by the particle pairs (the electron-positron $e^{\pm}$ and nucleon-antinucleon $N\bar{N}$ plasma) that have remained, in small quantities, from the epoch of annihilation, which occurred at red shifts

$$z_{ann} = \begin{cases} 2\cdot 10^9, & e^{\pm}, \\ 4\cdot 10^{12}, & N\bar{N}. \end{cases} \tag{4.8}$$

Calculation of the residual "quenched" concentration of e^+e^- and $N\bar{N}$ pairs in a charge-symmetric expanding universe at $z < z_{\rm ann}$ gives[7,19] a very small value

$$n = qn_\gamma, \quad q = \begin{cases} 10^{-16}, & e^{\pm}, \\ 3\cdot 10^{-19}, & N\bar{N}, \end{cases} \tag{4.9}$$

where $n_\gamma = 20T^3 \approx 400(1+z)^3\,{\rm cm}^{-3}$ is the photon concentration in the relic radiation and the ratio $q = n/n_\gamma$ is a quantity that is conserved in the expansion of the universe. The detection of such a negligible number of pairs of "quenched" matter from its electromagnetic annihilation radiation, the spectrum of which was analyzed in Ref. 19, is extremely unlikely.

We shall give an upper estimate for this radiation. We shall assume that, despite the extreme rarefaction of the matter, all the "quenched" matter is completely annihilated in the epoch of red shift $z_r \approx z_{\rm bound} \approx 4$. The energy density $\mathscr{E}_a$ of the annihilation radiation that was formed at z_r falls as $(1+z)^{-4}$ during the subsequent expansion of the universe (just as the relic-radiation energy $\mathscr{E}_\gamma = 2.7kTn_\gamma = \mathscr{E}_{\gamma 0}(1+z)^{-4}$), and, when recalculated for the observation time $z = 0$, amounts to (for $e^{\pm}$ and $N\bar{N}$ pairs)

$$\mathscr{E}_{a0} = 2mc^2\eta\frac{n(z_r)}{(1+z_r)^4} = 0.74\frac{\mathscr{E}_{\gamma 0}}{1+z_r}[\eta q z_{ann}] \approx 10^{-8}\ {\rm eV/cm^3}, \tag{4.10}$$

where $\mathscr{E}_{\gamma 0} = 0.25\ {\rm eV/cm^3}$ is the energy density of the relic radiation with temperature $T_0 = 2.7$ K, and η is a factor taking account of the processing of pair mass into electromagnetic radiation ($\eta_{e^{\pm}} \simeq 1$, $\eta_{N\bar{N}} \simeq 0.3$).

The energy of the annihilation photons, which at time z_r amounts to $\bar{\varepsilon}_{e^{\pm}} \approx 0.5$ MeV and $\bar{\varepsilon}_{N\bar{N}} \approx 200$ MeV, falls, by the time $z = 0$, by a factor of $1 + z_r \approx 5$, so that the annihilation radiation has energies in the region of 0.1 MeV and 40 MeV, respectively. The observed background radiation in these regions of energy has energy densities $\mathscr{E}(0.1\ {\rm MeV}) \approx 10^{-5}\ {\rm eV/cm^3}$ and $\mathscr{E}(40\ {\rm MeV}) \approx 10^{-7}\ {\rm eV/cm^3}$, which are much greater than the expected value (4.10). The observable background, in the indicated energy ranges, can be explained by the annihilation of pairs if the concentration of antiparticles outside the island is several orders greater than (4.9). For example, if there is explicit as well as spontaneous *CP*-violation, a situation is possible in which our baryon island is situated in a sea of antimatter, the concentration of which is only a few times lower than the concentration of matter in the island. In this case the background radiation could be significant.

C. Observations of individual sources

It is evident that none of the evolution processes inside a homogeneous island will give us any valuable information about its size or about our position with respect to the center, since information from different directions arrives at the observer simultaneously, with preservation of the spherical symmetry. Therefore, there is no hope of exhibiting any part of the celestial distribution of the discrete sources that are observable in different energy ranges if there is no certainty that they are being detected out to the edge of the island. From observations of extragalactic radio-sources and determinations of large numbers of their red shifts (from optical observations) it follows that a large part of the data corresponds to $z < 1$. It is evident that the same also applies to multiple observations in the infrared and x-ray regions. Therefore, a direct search for objects with large red shifts in different directions of the celestial sphere appears to be the most promising. The limiting value of z will give the size of the island, and the celestial distribution of the sources will make it possible to obtain the position of the observer with respect to the center. The most suitable objects may be quasars, young galaxies, clusters of galaxies, and clouds of intergalactic gas.

On Fig. 5 we show the position of the observer with respect to the center of the island and the event horizon, both of which he can see. The point C is the center of the island, r_C is the radius of the island, the point O is the position of the observer, $l_{\rm max}$ and $l_{\rm min}$ are the maximum and minimum metric distances from the observer to the edge of the island, and Δ is the displacement of the observer with respect to the center of the island. The angle φ is measured from the direction of the most distant point on the edge.

We shall denote by l the present-day metric distance to the edge of the island. In the Einstein–deSitter model, $l = 2cH_0^{-1}(1 - X^{1/2})$, where $X = (1+z)^{-1}$, z being the red shift of an object situated at the point A. From the trian-

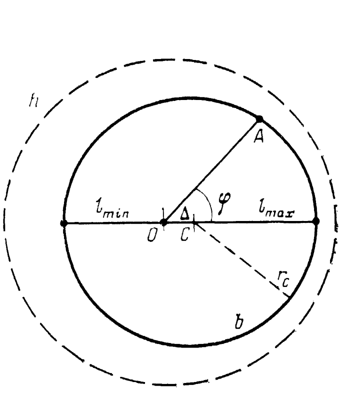

FIG. 5. Schematic arrangement of the island boundary (b) and the horizon (h).

gle AOC we find the celestial distribution of the boundary value of the red shift:

$$X=\{1-\beta[\varkappa\cos\varphi+(1-\varkappa^2\sin^2\varphi)^{1/2}]\}^2, \quad (4.11)$$

where $\varkappa$ is the ratio of the distance of the observer from the center of the island to the radius of the island:

$$\varkappa=\frac{\Delta}{r_C}=\frac{X_{max}^{1/2}-X_{min}^{1/2}}{2-X_{max}^{1/2}-X_{min}^{1/2}}, \quad (4.12)$$

and β is the ratio of the radius of the island to the horizon radius $h=2c/H_0$:

$$\beta=\frac{r_C H_0}{2c}=\frac{1}{2}(2-X_{max}^{1/2}-X_{min}^{1/2})\simeq 0.6. \quad (4.13)$$

The red shift for objects situated at the center of the island is described by the expression

$$X_C=(1+z_C)^{-1}=(1-\varkappa\beta)^2. \quad (4.14)$$

Above it was shown that $z_* - z_{**} < 0.2$; this is a bound on the eccentricity of our position in the island, and corresponds to the bounds $\varkappa < 0.01$ and $z_C < 0.01$.

CONCLUSION

The estimates given show that the model of a baryon island does not contradict the existing observations, although it does impose stringent restrictions on the position of the observer inside the island.

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[1)]Energetic quanta with energy $E \geqslant m_e c^2/6 \approx 80$ keV are scattered with higher probability by photons of the relic radiation than by "quenched" e^+e^- pairs.[19] For the energetic quanta the universe becomes transparent ($\tau_{\gamma\gamma} < 1$) at red shifts $z \ll z_{\gamma\gamma} \approx 2.5 \cdot 10^6 \, (m_e c^2/E)^{3/4}$.

[1]W. H. Sorrell, Mon. Not. R. Astron. Soc. **213**, 389 (1985); D. Crampton, A. P. Cowley, and F. D. A. Hartwick, Astrophys. J. **314**, 129 (1987).

[2]D. V. Korol'kov and Yu. N. Pariĭskiĭ, Soobshch. Shemakh. Astrofiz. Obs. Akad. Nauk Az. SSR No. 41, 42, 43 (1985).

[3]J. M. Uson and D. T. Wilkinson, Astrophys. J. Lett. **277**, L1 (1984).

[4]N. A. Zabotin and P. D. Nasel'skiĭ, Astron. Zh. **59**, 447 (1982); **60**, 430 (1983) [Sov. Astron. **26**, 272 (1982); **27**, 251 (1983)].

[5]N. Kardaschev, H. Blome, and W. Priester, Preprint, Space Research Institute, Moscow (1987).

[6]A. D. Dolgov and N. S. Kardaschev, Preprint No. 11 p-1190 (1986); Nature (in press).

[7]Ya. B. Zel'dovich and I. D. Novikov, *Relativistic Astrophysics, Vol. 2. The Structure and Evolution of the Universe*, Univ. of Chicago Press, Chicago (1983) [Russ. original, Nauka, Moscow (1975)].

[8]T. D. Lee, Phys. Rev. D **8**, 1226 (1973).

[9]D. A. Kirzhnits, Pis'ma Zh. Eksp. Teor. Fiz. **15**, 745 (1972) [JETP Lett. **15**, 529 (1972)]; D. A. Kirzhnits and A. D. Linde, Phys. Lett. **42B**, 471 (1972).

[10]S. Coleman and E. Weinberg, Phys. Rev. D **7**, 1888 (1983).

[11]A. D. Linde, Phys. Lett. **108B**, 389 (1982); A. Albrecht and P. J. Steinhardt, Phys. Rev. Lett. **48**, 1220 (1982).

[12]L. F. Abbott, E. Farhi, and M. B. Wise, Phys. Lett. **117B**, 29 (1982); A. Albrecht, P. J. Steinhardt, M. S. Turner, and F. Wilczek, Phys. Rev. Lett. **48**, 1437 (1982); A. D. Dolgov and A. D. Linde, Phys. Lett. **116B**, 329 (1982).

[13]A. H. Guth and S.-Y. Pi, Phys. Rev. Lett. **49**, 1110 (1982); S. W. Hawking, Phys. Lett. **115B**, 295 (1982); A. A. Starobinsky, Phys. Lett. **117B**, 175 (1982); J. M. Bardeen, P. J. Steinhardt, and M. S. Turner, Phys. Rev. D **28**, 679 (1983).

[14]V. A. Kuzmin, M. E. Shaposhnikov, and I. I. Tkachev, Phys. Lett. **105B**, 167 (1982).

[15]Ya. B. Zel'dovich, I. Yu. Kobzarev, and L. B. Okun', Zh. Eksp. Teor. Fiz. **67**, 3 (1974) [Sov. Phys. JETP **40**, 1 (1974)].

[16]R. K. Sachs and A. M. Wolfe, Astrophys. J. **147**, 73 (1967).

[17]J. Silk, Astrophys. J. **151**, 459 (1968).

[18]R. A. Sunyaev and Ya. B. Zel'dovich, Astrophys. Space Sci. **6**, 358 [**7**, No. 1, 3] (1968).

[19]A. F. Illarionov, Preprint No. 1273, Space Research Institute, Academy of Sciences of the USSR (1987).

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