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AN OBSTACLE TO CREATING A UNIVERSE IN THE LABORATORY*

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ABSTRACT

We show that any spherically symmetric false vacuum bubble which forms in an asymptotically flat space and grows beyond a certain critical size must have emerged from an initial singularity. Our result requires that the energy-momentum tensor obey the condition $T_{\mu\nu}k^\mu k^\nu \geq 0$ for all null k^μ , a property which holds at the classical level for almost all theories of matter. For the nonspherical case, we state a necessary condition that a false vacuum bubble must meet in order to avoid an initial singularity. We do not know if this condition can ever be met. The requirement of an initial singularity appears to be an insurmountable obstacle to the creation of an inflationary universe in the laboratory.

I. INTRODUCTION

The universe which we observe today was once much smaller. According to the inflationary model, the observed universe grew out of a region with a spatial size of less than 10^{-24} cm and a mass of less than 10 kg. The interior of this region consisted of metastable "false vacuum", characterized by a positive energy density and a negative pressure. The region grew exponentially with time for about 10^{-30} sec or more, until the false vacuum decayed. Since the inflationary epoch the universe has continued to expand, evolving in a manner well-described by standard big bang cosmology.

Since the inflationary model makes it possible for the universe to evolve from such a modest embryo, it is natural to ask whether we can create conditions in a small region of space which could give rise to a new universe. Suppose that a finite region of space, say the inside of a cauldron, has been arranged to be in a false vacuum state. Can this false vacuum region undergo exponential expansion to become a "child universe" with a future evolution similar to that of our own universe?

Unfortunately, the formation of a false vacuum requires an energy density that is far too high to be provided by any known technology. However, one may still ask whether it is in principle possible to create a universe by man-made processes. We find this question intriguing, and the analysis is possibly relevant to the early history of our universe.

The evolution of a region of false vacuum has been calculated [1-4] for an idealized system: a spherical bubble of false vacuum separated by a thin wall from an

Fig. 1. A spacetime diagram of a sample false vacuum bubble solution. Angular coordinates are suppressed, and lightlike lines travel at 45° . The bubble wall is shown as a heavy line with an arrow on it. The true vacuum region is to the right of the wall, where the metric describes part of Schwarzschild space, shown in Kruskal coordinates. The false vacuum region is to the left, where the metric describes part of de Sitter space. It is shown using peculiar coordinates chosen to join smoothly to the Schwarzschild region. The lower wiggly line indicates the initial singularity, and the upper wiggly line indicates the final singularity.

infinite region of true vacuum ($T_{\mu\nu} = 0$). In the true vacuum region the geometry is Schwarzschild, with the mass parameter determined by the combined energy of the bubble and the wall. Inside the bubble the geometry is de Sitter, with the cosmological constant determined by the false vacuum energy density. The equations of motion for the wall are found by requiring that the Einstein equations hold at the boundary [5].

The solutions of these equations have various properties depending on the initial conditions, but we will focus on those solutions for which the bubble grows without bound. An example of such a solution is shown in Fig. 1.

Although the figure does not make it evident, the proper volume of the de Sitter region and the proper area of the bubble wall both grow rapidly as time (i.e., the vertical coordinate of the diagram) increases. Nonetheless the expanding new universe does not encroach upon the "outside" true vacuum region. The non-Euclidean nature of the geometry allows the child universe to expand at no cost to the parent universe, and thus we would not be destroyed by a universe that we might create. To an observer in the true vacuum region the bubble wall is always inside the Schwarzschild horizon, and the newly created universe appears as a black hole. Only the earliest stages of the bubble evolution are in principle observable.

Taking these known solutions as a guide, it seems that the creation of a universe

is necessarily associated with the production of a black hole. This, however, does not, in principle present an insurmountable obstacle. Ordinary materials (e.g. stars) can collapse to form black holes, and it is possible at least to conceive of a laboratory setup that would produce a black hole.

However, there is one feature of the known solutions that must be avoided if we are to imagine creating a child universe—all of the solutions for which the bubble grows without limit have the property that the bubble is associated with an initial singularity, as in Fig. 1. This singularity is a spacelike boundary to the manifold, where any past directed geodesic which intersects it terminates. Such a singularity cannot conceivably be produced in the laboratory, since it has no prior history whatever. In the standard big bang model one hypothesizes that the universe originated from such an initial singularity, but we do not know of any initial singularities available in the universe today. (The black hole created in association with the child universe also has a final singularity—i.e., a boundary resulting in the termination of future directed geodesics. This singularity does not trouble us, however, since singularities of this type can be created by the collapse of ordinary matter.)

We would therefore like to know whether it is possible to avoid the initial singularity. By introducing some kind of "machinery" to complicate the evolution, can we start from nonsingular initial data and grow an arbitrarily large child universe? Unfortunately, the answer seems to be no. It appears that the outward velocity required for unlimited growth is so large that such a bubble can originate only from an initial singularity. In the next section we show that any child universe which grows beyond a certain size and approaches spherical symmetry must be associated with an initial singularity. Our proof uses a singularity theorem of general relativity first derived by Penrose. In section III we study the much more difficult case of configurations without spherical symmetry. Here we cannot be decisive, but we state a necessary condition that a geometry must satisfy in order to avoid an initial singularity. This condition seems rather difficult to satisfy, and even if it is met it does not guarantee the absence of a singularity. Finally, we discuss our conclusions in section IV.

II. SPHERICALLY SYMMETRIC CHILD UNIVERSES

Throughout this paper we assume that the interior of our "cauldron" is precisely in a false vacuum state, which means that the energy-momentum tensor is given by

$$T_{\mu\nu} = -\rho_f g_{\mu\nu}, \quad (1)$$

where ρ_f denotes the energy density of the false vacuum, assumed to be positive. The Einstein equations then take the form

$$R_{\mu\nu} = 3\chi^2 g_{\mu\nu}, \quad (2)$$

where

$$\chi^2 = \frac{8\pi}{3} G\rho_f. \quad (3)$$

In this section we also assume that the false vacuum region is spherically symmetric. We do not require spherical symmetry for the parent universe, however, provided that

the spherical symmetry of the false vacuum region is not perturbed. The analogue of Birkhoff's theorem then implies that the geometry for the false vacuum region is a part of de Sitter space.

To show that in this case the initial singularity is unavoidable, we begin by showing that de Sitter space contains 2-spheres that are "anti-trapped". We call a compact two dimensional surface anti-trapped if it has the property that both the inward and outward going sets of light rays, emanating normally from the surface, are diverging at every point on the surface. (If both sets of rays are converging at every point, the surface is called "trapped".) This result can be understood intuitively by viewing de Sitter space as an expanding space, with geodesics tending to move apart. It is natural, say in flat space, for the outward going light rays normal to a 2-sphere to be diverging, and for the inward going light rays to be converging. In an expanding space, however, the inward going light rays are pulled outward by the expansion, and this expansion can overcome the inward motion and lead to divergence if the 2-sphere is sufficiently large. This is what happens in de Sitter space.

We will describe de Sitter space, as well as other geometries, by using coordinates which are tailored to our purposes.

To construct the coordinates, choose a spacetime point P and in the neighborhood of P construct locally flat conventional spherical coordinates (t, r, θ, φ) , where $t = r = 0$ is the point P . For each (θ, φ) there is a future directed null geodesic leaving P headed in the (θ, φ) direction. Scale the affine parameter of each geodesic so that it equals r in the neighborhood of P , and continue to call the affine parameter r as the geodesics extend beyond the neighborhood of P . The coordinates (r, θ, φ) now label a three dimensional subspace which goes beyond the neighborhood of P as far as the geodesics can be followed. For fixed r the surface labelled by (θ, φ) is a 2-surface. From each point on this 2-surface, there is an inward directed null geodesic emanating normally from the 2-surface. Call the affine parameter along each of these geodesics v , with $v = 0$ being the point of departure from the surface. There is still an arbitrariness in the scale of v , which we fix by requiring that $n \cdot \ell = -1$ at $v = 0$, where $n^\mu \equiv \partial x^\mu / \partial r$ is the tangent vector to the outward directed null geodesic $x^\mu(r)$, and $\ell^\mu \equiv \partial x^\mu / \partial v$ is the tangent vector to the inward directed null geodesic $x^\mu(v)$. This construction gives good coordinates at least for v near zero and as long as the geodesics parameterized by r (at $v = 0$) can be followed and do not cross.

For the special case of de Sitter space, the metric in the coordinates (r, v, θ, φ) is found to be

$$ds^2 = -2dr dv + \left[r - \frac{v}{2} (1 - \chi^2 r^2) \right]^2 (d\theta^2 + \sin^2 \theta d\varphi^2) . \quad (4)$$

This metric can be found by starting with any familiar de Sitter space coordinates and metric and then explicitly solving for the null geodesics described in the previous paragraph. Alternatively, it suffices to check that the metric (4) satisfies (2) as well as the stated coordinate conditions.

Given the metric, it is easy to calculate the divergences of the inward and outward going null geodesics leaving the 2-sphere $r = \text{constant}$, $v = 0$. The four-velocity vectors for these geodesics are given respectively by $\ell^\mu = (0, 1, 0, 0)$ and

$n^\mu = (1, 0, 0, 0)$, so one has

$$\theta_{\text{in}} \equiv \nabla_\mu \ell^\mu \Big|_{v=0} = \Gamma_{\mu\nu}^\mu \Big|_{v=0} = -\frac{1}{r} (1 - r^2 \chi^2) \quad (5)$$

and

$$\theta_{\text{out}} \equiv \nabla_\mu n^\mu \Big|_{v=0} = \Gamma_{\mu\nu}^\mu \Big|_{v=0} = \frac{2}{r} . \quad (6)$$

We see that the inward going light rays become diverging for $r > 1/\chi$. Thus, for $r > 1/\chi$ the 2-spheres of fixed r are anti-trapped.

We can now use the Penrose theorem [7] (or see for example ref. [6], p. 239) to tell us when the existence of an anti-trapped surface necessarily implies an initial singularity. This theorem states that in a connected spacetime, if

- there exists a noncompact Cauchy surface,
- $R_{\mu\nu} k^\mu k^\nu \geq 0$ for all null k^μ , and
- there exists an anti-trapped surface,

then there exists an initial singularity. For the purposes of this theorem, an "initial singularity" is defined as the existence of at least one past directed null geodesic which cannot be extended beyond some finite range of its affine parameter. (For a precise definition of "inextendible", see ref. [6], p. 193. In detail, the theorem states that at least one past directed null geodesic cannot be extended beyond an affine length of $2/\theta_{\text{min}}$, where θ_{min} is the minimum value of the divergence of the inward and outward going null geodesics.) We stated the theorem in an unconventional form, using anti-trapped surfaces and initial singularities, while the usual statement has trapped surfaces and final singularities. Our statement, however, is easily seen to be the time reversal of the usual statement, since a future directed diverging set of null geodesics is also a past directed converging set.

We assume that the laboratory which is to produce the false vacuum bubble exists in a space which is asymptotically flat. If we further assume that the space has a well-defined Cauchy development, then the Cauchy surfaces are noncompact and condition (a) is satisfied. The relevance of this assumption to our own universe will be discussed in the last section.

As a consequence of the Einstein equations, condition (b) can be cast as a condition on the energy momentum tensor: $T_{\mu\nu} k^\mu k^\nu \geq 0$ for all null k^μ . We call this the "very weak energy condition" since it is less restrictive than the usual weak energy condition, $T_{\mu\nu} t^\mu t^\nu \geq 0$ for all timelike t^μ . Inside the bubble, where $T_{\mu\nu}$ is proportional to $g_{\mu\nu}$, the very weak energy condition is clearly satisfied. Outside the bubble, it is a condition on the matter distribution which says that the energy density plus any of the principal pressures is greater than or equal to zero. Any standard theory of matter in its classical approximation obeys the weak energy condition, and hence also the very weak condition. (The only familiar type of theory that leads to a violation of the very weak energy condition at the classical level is that of a scalar field coupled nonminimally to gravity, that is, with an $R\phi^2$ term in the Lagrangian.) We therefore make the

reasonable assumption that the machinery to be used in the construction of the child universe will obey the very weak energy condition, thereby satisfying condition (b).

For a spherically symmetric child universe we can describe the bubble interior by the metric (4), placing $r = 0$ at the center of the bubble. If the bubble is large enough so that this metric is valid for some r larger than $1/\chi$, then an anti-trapped surface exists. In this case an initial singularity is unavoidable.

Thus, singular initial conditions are required to produce any child universe which at some late time approaches spherical symmetry or has a spherically symmetric region which can be described by the metric (4) for some $r > 1/\chi$. (Again, we assume that the very weak energy condition holds and that the parent universe is asymptotically flat.) Since small perturbations away from a de Sitter space metric tend to diminish at large times [8,9], it is doubtful that a small deviation from spherical symmetry will be enough to avoid the necessity of singular initial conditions.

Finally, we must discuss whether a singularity as defined by the Penrose theorem poses an insurmountable obstacle to the creation of an inflationary universe. The situation was simpler for the known exact solutions, since in those cases the $r = 0$ hypersurface could be taken as a spacelike boundary to the manifold, providing an initial hypersurface for which there was no prior history. When one is relying on the Penrose theorem, on the other hand, one knows only that there exists at least one null geodesic which cannot be continued arbitrarily far into the past. This means, however, that the history of this geodesic is incomplete, and the presence or absence of a light ray traveling along the geodesic cannot be determined by initial conditions set prior to the earliest time along the geodesic. We conclude, therefore, that such a situation cannot arise by well-behaved Cauchy development from smooth initial conditions, and thus it seems reasonable to regard this kind of a singularity as an insurmountable obstacle.

III. CHILD UNIVERSES WITHOUT SPHERICAL SYMMETRY

In the spherically symmetric case we constructed 2-spheres by following, out to a fixed affine parameter r , the light rays emanating in all directions from a spacetime point P . We could show that for large enough r these 2-spheres are anti-trapped surfaces. To do so we used the fact that the metric is that of de Sitter space, which for our choice of coordinates has the form (4).

An arbitrary false vacuum bubble will satisfy the Einstein equations (2), but these equations do not determine the metric without further information. The matter distribution outside the bubble affects the metric inside the bubble, whereas the local conditions (2) are maintained as long as the false vacuum region is empty.

Nonetheless, it is tempting to conjecture that if a large enough region of spacetime obeys eq. (2), then it will contain an anti-trapped surface. Unfortunately, we have not been able to establish this result. However, we can show that if eq. (2) holds in a sufficiently large region, then one can construct a 2-surface for which the average of the divergence of the inward directed null geodesics normal to the surface is positive.

To derive this result we again look at 2-surfaces defined by following the light rays emanating from some fixed point P out to a fixed value of the affine parameter. To do so it is convenient to establish the coordinates $(\tau, v, \theta, \varphi)$ by following the prescription

given in the last section. This prescription leads to a coordinate system in which the metric has certain properties. Since $\ell^\mu = (0, 1, 0, 0)$ is a geodesic, $\ell^\alpha \nabla_\alpha \ell^\mu = 0$; and since r and v are coordinates, $n^\alpha \nabla_\alpha \ell^\mu = \ell^\alpha \nabla_\alpha n^\mu$, where $n^\mu = (1, 0, 0, 0)$. Together, these imply that $\frac{\partial}{\partial v} g_{rv} = \ell^\mu \nabla_\mu (n \cdot \ell) = 0$; since $g_{rv} = -1$ at $v = 0$, it follows that $g_{rv} = -1$ everywhere. By similar arguments one can show that $g_{vv} = g_{v\theta} = g_{v\varphi} = 0$ everywhere. The metric in the $(\tau, v, \theta, \varphi)$ coordinates has the form:

$$g_{\mu\nu} = \begin{bmatrix} g_{\tau\tau} & -1 & g_{r\theta} & g_{r\varphi} \\ -1 & 0 & 0 & 0 \\ g_{r\theta} & 0 & g_{\theta\theta} & g_{\theta\varphi} \\ g_{r\varphi} & 0 & g_{\theta\varphi} & g_{\varphi\varphi} \end{bmatrix}. \quad (7)$$

We can use the fact that n is a geodesic at $v = 0$ to show that $\frac{\partial}{\partial r} g_{r\theta}(v = 0) = 0$. Since for r small and $v = 0$ the metric takes its flat space form with $g_{r\theta} = 0$, we see that

$$g_{r\theta}(v = 0) = 0, \quad (8a)$$

and similarly

$$g_{r\varphi}(v = 0) = 0. \quad (8b)$$

Also, note that eq. (7) implies that the determinant $|g| = g^{(2)}$, where $g^{(2)}$ is the determinant of the two-dimensional metric g_{ij} with $(i, j) = (\theta, \varphi)$.

Now look at the $(i, j) = (\theta, \varphi)$ component of eq. (2). The Ricci tensor R_{ij} can be decomposed into the two dimensional Ricci tensor $R_{ij}^{(2)}$ (constructed from the two dimensional metric g_{ij}) and a remainder R_{ij}' . (It is useful to note that, given the metric (7), g^{ij} depends only on g_{ij} .) Contracting eq. (2) with g^{ij} gives

$$6\chi^2 = R^{(2)} + g^{ij} R_{ij}', \quad (9)$$

where $R^{(2)}$ is the two-dimensional curvature scalar. Evaluating R_{ij}' at $v = 0$ gives

$$\begin{aligned} \frac{1}{\sqrt{g^{(2)}}} \frac{\partial}{\partial \tau} \frac{\partial}{\partial v} \sqrt{g^{(2)}} &= 3\chi^2 - \frac{R^{(2)}}{2} + \frac{1}{2\sqrt{g^{(2)}}} \partial_j \left(g^{ij} g_{ir, v} \sqrt{g^{(2)}} \right) \\ &+ \frac{1}{4} g^{ij} g_{ir, v} g_{j\tau, v} \quad (v = 0). \end{aligned} \quad (10)$$

Following Hartle and Wilkins [10], we integrate this equation over the 2-surface $r = \text{constant}$, $v = 0$. First, the integral of the scalar curvature $R^{(2)}$ over the 2-surface is 8π by the Gauss-Bonnet theorem. The next term is a total divergence and vanishes upon integration. The last term is positive so we get the inequality

$$\frac{\partial}{\partial \tau} \frac{\partial}{\partial v} a(\tau, v) \Big|_{v=0} \geq 3\chi^2 a(\tau, v = 0) - 4\pi, \quad (11)$$

where $a(r, v)$ is the area of the 2-surface. Integrating from 0 out to r gives

$$\frac{\partial}{\partial v} a(r, v) \Big|_{v=0} \geq 3\chi^2 \int_0^r a(r, v=0) dr - 4\pi r, \quad (12)$$

where we have used the fact that $a(r, v)$ approaches its flat space value of zero as r and v approach zero.

Now $\theta_{1n} = \nabla_\mu \ell^\mu \Big|_{v=0}$ which for our coordinates is

$$\theta_{1n} = \Gamma_{\mu\nu}^\mu \Big|_{v=0} = \frac{1}{\sqrt{g^{(2)}}} \frac{\partial}{\partial v} \sqrt{g^{(2)}} \Big|_{v=0}. \quad (13)$$

If we average θ_{1n} over the 2-surface we obtain

$$\langle \theta_{1n} \rangle \equiv \frac{\int_{\text{surface}} \sqrt{g^{(2)}} \theta_{1n}}{\int_{\text{surface}} \sqrt{g^{(2)}}} = \frac{1}{a(r, v)} \frac{\partial}{\partial v} a(r, v) \Big|_{v=0}, \quad (14)$$

which combines with (12) to give

$$\langle \theta_{1n} \rangle \geq \frac{3\chi^2 r}{a(r, v=0)} \left[\frac{1}{r} \int_0^r a(r, v=0) dr - \frac{4\pi}{3\chi^2} \right]. \quad (15)$$

Imagine following the light rays emanating from some point P out to affine parameter r , and for each r compute the area $a(r)$ of the 2-surface of fixed r (for $v=0$). We assume that $a(r)$ increases since we are following outward going light rays in an expanding space. When the value of $a(r)$ averaged over r , i.e. $\frac{1}{r} \int_0^r a(r) dr$, exceeds $4\pi/3\chi^2$, we know that $\langle \theta_{1n} \rangle$, the average of θ_{1n} over the 2-surface of fixed r , is positive. For de Sitter space this happens at $r = 1/\chi$ and the inequality (15) is actually an equality.

In order to apply the Penrose theorem (in its anti-trapped form) we need to know that θ_{1n} is positive everywhere on a 2-surface—it is not enough for the average to be positive. (We assume that θ_{out} is positive everywhere on the 2-surface, since outward moving null geodesics are naturally diverging.) However, to show that an initial singularity exists we need to find only one anti-trapped surface. In order to avoid an initial singularity, the false vacuum region must have the property that for each point P , all of the associated 2-surfaces of constant r have some point where θ_{1n} is negative, even though the average of θ_{1n} over the 2-surface is positive, for r large enough. This property still does not guarantee the absence of an initial singularity; it is a necessary condition only. We have not succeeded in finding a geometry which has this property.

IV. DISCUSSION

As you may have imagined, it is rather difficult to create a new universe in the laboratory. Assuming that the very weak energy condition holds, then any spherically symmetric child universe that is created in a asymptotically flat parent universe is

associated with an initial singularity. The same appears to be true for child universes without spherical symmetry, although we have not excluded the possibility that the singularity might be avoided by a sufficiently weird bubble geometry.

Since the big-bang description of our universe begins with an initial singularity, we should ask whether the creation of a child universe requires a new singularity in addition to the big bang. In addition, we should ask whether our results can be applied if our universe has compact Cauchy surfaces, as would be the case if the universe were closed. The answer to both questions appears to be yes, for similar reasons. We imagine that the laboratory which produces the false vacuum bubble is very small compared to the Hubble length, and is to accomplish its mission in a time which is very short compared to the Hubble time. The cosmological spacetime curvature is then very small on the relevant scales of length and time. Although a high energy density is obviously necessary for the creation of a false vacuum bubble, we nonetheless assume that the energy used to construct and operate the laboratory was initially very dilute—the high density is obtained by compressing a modest amount of energy into an incredibly small volume. Thus, prior to this compression the perturbation in the spacetime metric was extremely small, and the spacetime was well-approximated, on the scale of the laboratory, by one which is Minkowskian both at large distances and at early times. Since the requirement of an initial singularity would preclude the creation of a child universe in such an asymptotically Minkowskian background, it seems reasonable to infer that the same conclusion would apply in the background provided by our own universe*.

Finally, our entire discussion has been in the context of classical general relativity. There is always the possibility that quantum mechanics could help us avoid a classical singularity. In particular, the quantum expectation value of the energy momentum tensor will not in general obey the very weak energy condition. We are presently investigating this and other routes to nonsingular, quantum universe creation.

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* However, if somewhere in our universe there exists a "white hole", a piece of the initial big bang singularity, then the creation of a child universe from such an object is not excluded by our arguments.

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