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Research Article

Could Some Black Holes Have Evolved from Wormholes?

Peter K. F. Kuhfittig

Department of Mathematics, Milwaukee School of Engineering, Milwaukee, WI 53202-3109, USA

Correspondence should be addressed to Peter K. F. Kuhfittig, kuhfitti@msoe.edu

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One way to explain the present acceleration of the universe is Einstein's cosmological constant. It is quite likely, in view of some recent studies, that a time-dependent equation of state had caused the Universe to evolve from an earlier phantom-energy model. In that case, traversable wormholes could have formed spontaneously. It is shown in this paper that such wormholes would eventually have become black holes. This would provide a possible explanation for the huge number of black holes discovered, while any evidence for the existence of wormholes is entirely missing, even though wormholes are just as good, in terms of being a prediction of general relativity, as black holes.

1. Introduction

Traversable wormholes, whose possible existence was first conjectured by Morris and Thorne in 1988 [1], are shortcuts or tunnels that could in principle be used for traveling to remote parts of our Universe or to different universes altogether. Using units in which $c = G = 1$, a wormhole can be described by the general line element

$$ds^2 = -e^{2\gamma(r)}dt^2 + e^{2\alpha(r)}dr^2 + r^2(d\theta^2 + \sin^2\theta\, d\phi^2). \tag{1}$$

The motivation for this idea is the Schwarzschild line element

$$ds^2 = -\left(1 - \frac{2M}{r}\right)dt^2 + \frac{dr^2}{1 - 2M/r} + r^2(d\theta^2 + \sin^2\theta\, d\phi^2), \tag{2}$$

which may be viewed as a black hole centered at the origin of a (t, r, θ, ϕ) coordinate system. Both wormholes and black holes are predictions of Einstein's general theory of relativity. The main difference is that wormholes must necessarily violate certain energy conditions; more precisely, the stress-energy tensor of the matter source of gravity violates the weak energy condition [1]. If the matter source was different in the past, then wormholes could theoretically have been formed naturally. Examples are equations of state that parameterize certain dark energy models, as we will see in the next section. Moreover, according to some recent studies, such a transition is likely to have occurred in the relatively recent past [2, 3].

In this paper, we study an equation of state that is both space and time dependent. It is proposed that if the equation of state evolved into the present cosmological constant model, then any wormhole that had previously come into existence would have formed an event horizon, thereby becoming a black hole. This would provide a possible explanation for the failure to detect any evidence of wormholes, while black holes appear to be abundant.

2. Background

Interest in traversable wormholes has increased in recent years due to an unexpected connection, the discovery that our Universe is undergoing an accelerated expansion [4, 5]. This acceleration, caused by a negative pressure *dark energy*, implies that $\ddot{a} > 0$ in the Friedmann equation $\ddot{a}/a = -(4\pi/3)(\rho + 3p)$. The equation of state is $p = -K\rho$, $K > 1/3$, and $\rho > 0$. While the condition $K > 1/3$ is required for an accelerated expansion, larger values for K are also of interest. In fact, the most appealing candidate for dark energy is $K = 1$, corresponding to the cosmological constant Λ [6]. The presence of the cosmological constant has resulted in a modification of the Einstein field equations by effectively adding an isotropic and homogeneous source with constant equation of state [7]. One can therefore argue, as in [8], that this model is the primary candidate for the present Universe.

Another widely studied possibility is the case $K > 1$, referred to as *phantom energy* [9]. To see why, we need to recall that the set of *orthonormal basis vectors* may be interpreted as the proper frame of a set of observers who remain at rest in the coordinate system. If the basis vectors are denoted by e_t, e_r, e_θ, and e_ϕ, then the orthonormal basis vectors are (referring to line element (1))

$$e_{\hat{t}} = e^{-\gamma(r)} e_t, \qquad e_{\hat{r}} = e^{-\alpha(r)} e_r,$$
$$e_{\hat{\theta}} = r^{-1} e_\theta, \qquad e_{\hat{\phi}} = (r \sin\theta)^{-1} e_\phi. \tag{3}$$

In this frame of reference, the components of the stress-energy tensor $T_{\hat{\alpha}\hat{\beta}}$ have an immediate physical interpretation: $T_{\hat{t}\hat{t}} = \rho$, $T_{\hat{r}\hat{r}} = p$, $T_{\hat{\theta}\hat{\theta}} = T_{\hat{\phi}\hat{\phi}} = p_t$, where ρ is the energy density, p the radial pressure, and p_t the lateral pressure. The weak energy condition (WEC) can now be stated as follows: $T_{\hat{\alpha}\hat{\beta}} \mu^{\hat{\alpha}} \mu^{\hat{\beta}} \geq 0$ for all time-like vectors and, by continuity, all null vectors. For example, given the radial outgoing null vector $(1, 1, 0, 0)$, we obtain

$$T_{\hat{\alpha}\hat{\beta}} \mu^{\hat{\alpha}} \mu^{\hat{\beta}} = \rho + p \geq 0. \tag{4}$$

So if the WEC is violated, we have $\rho + p < 0$. While all classical forms of matter ordinarily meet this condition, there are situations in quantum field theory, such as the Casimir effect, that allow such violations [10]. More importantly, in our case this violation occurs whenever $K > 1$, thereby meeting the primary prerequisite for the existence of wormholes.

Two recent papers [11, 12] discuss wormhole solutions that depend on a variable equation of state parameter, that is, $p/\rho = -K(r)$, where $K(r) > 1$ for all r. The variable r refers to the radial coordinate in the line element

$$ds^2 = -e^{2h(r)} dt^2 + \frac{1}{1 - b(r)/r} dr^2 + r^2 (d\theta^2 + \sin^2\theta \, d\phi^2). \tag{5}$$

In this form of the line element, $h = h(r)$ is called the *redshift function* and $b = b(r)$ is the *shape function*. The minimum radius $r = r_0$ corresponds to the *throat* of the wormhole, where $b(r_0) = r_0$.

It is shown in [12] that given a specific shape function, it is possible to determine $K(r)$ and vice versa. It is also assumed that $h'(r) \equiv 0$, referred to as the "zero tidal-force solution" in [1].

An earlier study [13] assumed that the equation of state is time dependent. In this paper, we will assume that the equation depends on both r and t, while remaining independent of direction.

Since the equation of state is time dependent, we assume that the corresponding metric is also time dependent. It is shown in the next section that such a metric describes a slowly evolving wormhole structure without assigning specific functions to h, b, and K. In particular, the function h in line element (5) need not be a constant.

Evolving wormhole geometries are also discussed in [14, 15].

Since we are dealing with a given time-dependent equation of state, it is natural to consider the consequences of an evolving equation, particularly one in which the parameter $K(r, t)$ approaches unity. That is the topic of Section 6.

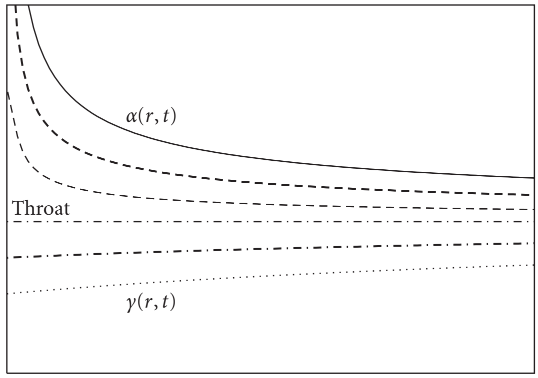

Figure 1: Graph showing the qualitative features of $\alpha(r, t)$ and $\gamma(r, t)$.

3. The Metric

In this paper, we will be dealing with a time-dependent metric describing an evolving wormhole

$$ds^2 = -e^{2\gamma(r,t)} dt^2 + e^{2\alpha(r,t)} dr^2 + r^2 (d\theta^2 + \sin^2\theta \, d\phi^2), \tag{6}$$

where γ and α have continuous partial derivatives, so that γ and α are continuous, as well.

In view of line element (5), we have

$$e^{2\alpha(r,t)} = \frac{1}{1 - b(r, t)/r}. \tag{7}$$

So the shape function is now given by

$$b(r, t) = r\left(1 - e^{-2\alpha(r,t)}\right). \tag{8}$$

To study the effect of a gradually changing equation of state, we assume the existence of a fixed throat at $r = r_0$, that is, $b(r_0, t) = r_0$ for all t. (In other words, the wormhole is close to being static for relatively long periods of time.) As a consequence, for all t,

$$\lim_{r \to r_0+} \alpha(r, t) = +\infty, \qquad \lim_{r \to r_0+} \frac{\partial}{\partial r} \alpha(r, t) = -\infty. \tag{9}$$

As before, $\gamma = \gamma(r, t)$ is the redshift function, which must be everywhere finite to prevent an event horizon at the throat. As in the case of the Schwarzschild line element, $(\partial/\partial r)\gamma(r, t) > 0$. The qualitative features of α and γ are shown in Figure 1.

To obtain a traversable wormhole, the shape function must obey the usual flare-out conditions at the throat, modified to accommodate the time dependence

$$b(r_0, t) = r_0, \quad \frac{\partial}{\partial r} b(r_0, t) < 1 \quad \forall t. \tag{10}$$

So $b(r, t) < r$ for all t near the throat. Another requirement is asymptotic flatness: $b(r, t)/r \to 0$ as $r \to \infty$.

The next step is to list the time-dependent components of the Einstein tensor in the orthonormal frame. (For a derivation, see Kuhfittig [16].)

$$G_{\hat{t}\hat{t}} = \frac{2}{r}e^{-2\alpha(r,t)}\frac{\partial}{\partial r}\alpha(r,t) + \frac{1}{r^2}\left(1 - e^{-2\alpha(r,t)}\right),$$

$$G_{\hat{r}\hat{r}} = \frac{2}{r}e^{-2\alpha(r,t)}\frac{\partial}{\partial r}\gamma(r,t) - \frac{1}{r^2}\left(1 - e^{-2\alpha(r,t)}\right), \tag{11}$$

$$G_{\hat{t}\hat{r}} = \frac{2}{r}e^{-\gamma(r,t)}e^{-\alpha(r,t)}\frac{\partial}{\partial t}\alpha(r,t),$$

$$\begin{aligned}G_{\hat{\theta}\hat{\theta}} = G_{\hat{\phi}\hat{\phi}} = &-e^{-2\gamma(r,t)} \\ &\times\left[\frac{\partial^2}{\partial t^2}\alpha(r,t) - \frac{\partial}{\partial t}\gamma(r,t)\frac{\partial}{\partial t}\alpha(r,t) + \left(\frac{\partial}{\partial t}\alpha(r,t)\right)^2\right] \\ &- e^{-2\alpha(r,t)}\left[-\frac{\partial^2}{\partial r^2}\gamma(r,t) + \frac{\partial}{\partial r}\gamma(r,t)\frac{\partial}{\partial r}\alpha(r,t) - \left(\frac{\partial}{\partial r}\gamma(r,t)\right)^2\right] \\ &- \frac{1}{r}e^{-2\alpha(r,t)}\left(-\frac{\partial}{\partial r}\gamma(r,t) + \frac{\partial}{\partial r}\alpha(r,t)\right).\end{aligned} \tag{12}$$

Recall that from the Einstein field equations in the orthonormal frame, $G_{\hat{\alpha}\hat{\beta}} = 8\pi T_{\hat{\alpha}\hat{\beta}}$, the components of the Einstein tensor are proportional to the components of the stress-energy tensor. In particular, $T_{\hat{t}\hat{t}} = \rho(r,t)$, $T_{\hat{r}\hat{r}} = p(r,t)$, $T_{\hat{\theta}\hat{\theta}} = T_{\hat{\phi}\hat{\phi}} = p_t(r,t)$, and $T_{\hat{t}\hat{r}} = T_{\hat{r}\hat{t}} = (1/8\pi)G_{\hat{t}\hat{r}} = g(r,t)$, where $f(r,t) = -g(r,t)$ is usually interpreted as the energy flux in the outward radial direction [17]. For the outgoing null vector $(1,1,0,0)$, the WEC, $T_{\hat{\alpha}\hat{\beta}}\mu^{\hat{\alpha}}\mu^{\hat{\beta}} \geq 0$, now becomes $\rho + p \pm 2g \geq 0$.

4. The Equation of State

Since (6) describes a slowly evolving wormhole, the equation of state should have a time-dependent parameter $K(r,t)$. As in [12], K depends on the radial coordinate, but not on the direction. To be compatible with the wormhole geometry in Section 3, Lobo suggested the inclusion of a term analogous to the flux term [18]. One such possibility is

$$p(r,t) = -K(r,t)[\rho(r,t) + 2g(r,t)], \tag{13}$$

where $|2g(r,t)| < \rho(r,t)$. The last condition implies that $\alpha(r,t)$ changes slowly enough; in fact, $g(r,t)$ can be identically zero. If $K(r,t) > 1$, this equation of state describes a generalized phantom-energy model, as we will see in (18). As a result, the case $K = 1$ would still correspond to a cosmological constant.

Since the notion of dark or phantom energy applies only to a homogeneous distribution of matter in the Universe, while wormhole spacetimes are necessarily inhomogeneous, we adopt the point of view in Sushkov [19]: extended to spherically symmetric wormhole geometries, the pressure appearing in the equation of state is now a negative radial pressure, while the transverse pressure is determined from the field equations.

Given the evolving equation of state (13), suppose at some point in the past, the equation of state parameter K had actually crossed the phantom divide, so that, for a time, $K(r,t) > 1$. In fact, according to [2, 3], it is quite likely that such a transition had taken place in the relatively recent past. Suppose further that K was decreasing, that is, $(\partial/\partial t)K(r,t) < 0$ with $K(r,t) \to 1$ at a time closer to the present, and that K decreased fast enough during this time interval to compensate for the very gradual increase in the energy density characteristic of phantom energy, allowing us to assume that

$$\frac{\partial}{\partial t}p(r,t) \geq 0. \tag{14}$$

With this information we can now determine the sign of $g(r,t)$:

$$\begin{aligned}\frac{\partial}{\partial t}p(r,t) &= \frac{1}{8\pi}\frac{\partial G_{\hat{r}\hat{r}}}{\partial t} \\ &= \frac{1}{8\pi}e^{-2\alpha(r,t)}\left[\frac{\partial}{\partial t}\alpha(r,t)\left((-2)\frac{2}{r}\frac{\partial}{\partial r}\gamma(r,t) - \frac{2}{r^2}\right) + \frac{2}{r}\frac{\partial}{\partial t}\left(\frac{\partial}{\partial r}\gamma(r,t)\right)\right] \geq 0.\end{aligned} \tag{15}$$

Since the pressure would change at a finite rate, it follows that $(\partial/\partial t)\alpha(r,t)$ is finite. Also,

$$\frac{\partial}{\partial t}\alpha(r,t) \leq 0 \quad \forall r, \tag{16}$$

so that $g(r,t) \leq 0$, as well. The reason is that, judging from Figure 1, the rate of change of the slope of γ, that is,

$$\frac{\partial}{\partial t}\left(\frac{\partial}{\partial r}\gamma(r,t)\right) \tag{17}$$

is likely to be vanishingly small. We will return to this point in Section 6.

Returning to (13), observe that

$$\begin{aligned}T_{\hat{\alpha}\hat{\beta}}\mu^{\hat{\alpha}}\mu^{\hat{\beta}} &= \rho + p + 2g \\ &= \rho(r,t) - K(r,t)[\rho(r,t) + 2g(r,t)] + 2g(r,t) \\ &= [\rho(r,t) + 2g(r,t)][-K(r,t) + 1] < 0,\end{aligned} \tag{18}$$

since $K(r,t) > 1$ and $|2g(r,t)| < \rho(r,t)$. So the WEC is violated, as one would expect in a phantom-energy scenario. We would also expect the flare-out conditions to be met. That is the topic of the next section.

5. The Flare-Out Conditions

From the Einstein field equations $G_{\hat{\alpha}\hat{\beta}} = 8\pi T_{\hat{\alpha}\hat{\beta}}$ and the above equation of state, we have

$$G_{\hat{t}\hat{t}} = 8\pi\rho, \qquad G_{\hat{r}\hat{r}} = 8\pi[-K(r,t)][\rho(r,t) + 2g(r,t)]. \tag{19}$$

Using (11), we obtain the following system of equations:

$$G_{\hat{t}\hat{t}} = 8\pi T_{\hat{t}\hat{t}} = \frac{2}{r}e^{-2\alpha(r,t)}\frac{\partial}{\partial r}\alpha(r,t) + \frac{1}{r^2}\left(1 - e^{-2\alpha(r,t)}\right),$$
$$G_{\hat{r}\hat{r}} = 8\pi T_{\hat{r}\hat{r}} = 8\pi[-K(r,t)][\rho(r,t) + 2g(r,t)] \quad (20)$$
$$= \frac{2}{r}e^{-2\alpha(r,t)}\frac{\partial}{\partial r}\gamma(r,t) - \frac{1}{r^2}\left(1 - e^{-2\alpha(r,t)}\right).$$

Substituting the expressions for $\rho(r,t)$ and $2g(r,t)$ yields

$$-K(r,t)\frac{2}{r}e^{-2\alpha(r,t)}\frac{\partial}{\partial r}\alpha(r,t)$$
$$= K(r,t)\frac{1}{r^2}\left(1 - e^{-2\alpha(r,t)}\right) + \frac{2}{r}e^{-2\alpha(r,t)}\frac{\partial}{\partial r}\gamma(r,t)$$
$$- \frac{1}{r^2}\left(1 - e^{-2\alpha(r,t)}\right) + K(r,t)\frac{4}{r}e^{-\gamma(r,t)}e^{-\alpha(r,t)}\frac{\partial}{\partial t}\alpha(r,t). \quad (21)$$

The following rearrangement will be needed again in Section 6:

$$\frac{\partial}{\partial r}\alpha(r,t) + 2e^{-\gamma(r,t)}e^{\alpha(r,t)}\frac{\partial}{\partial t}\alpha(r,t)$$
$$= -\frac{1}{2r}\left(e^{2\alpha(r,t)} - 1\right) + \frac{1}{2r}\frac{1}{K(r,t)}\left(e^{2\alpha(r,t)} - 1\right) \quad (22)$$
$$- \frac{1}{K(r,t)}\frac{\partial}{\partial r}\gamma(r,t).$$

For the purpose of analysis, however, a more convenient form is the following:

$$\frac{2(\partial/\partial r)\alpha(r,t)}{e^{2\alpha(r,t)} - 1} = \frac{-2(\partial/\partial r)\gamma(r,t)}{K(r,t)(e^{2\alpha(r,t)} - 1)} - \frac{1}{r}\left(1 - \frac{1}{K(r,t)}\right)$$
$$- \frac{4e^{-\gamma(r,t)}e^{\alpha(r,t)}(\partial/\partial t)\alpha(r,t)}{e^{2\alpha(r,t)} - 1}. \quad (23)$$

At this point, we define

$$F(r,t) = \frac{-2(\partial/\partial r)\gamma(r,t)}{K(r,t)(e^{2\alpha(r,t)} - 1)} - \frac{1}{r}\left(1 - \frac{1}{K(r,t)}\right)$$
$$- \frac{4e^{-\gamma(r,t)}e^{\alpha(r,t)}(\partial/\partial t)\alpha(r,t)}{e^{2\alpha(r,t)} - 1}. \quad (24)$$

The form $\int du/(e^u - 1) = \ln(e^u - 1) - u$ now yields

$$\ln\left(e^{2\alpha(r,t)} - 1\right) - 2\alpha(r,t) = \int_c^r F(r',t)dr', \quad (25)$$

where c is an arbitrary constant. Recalling that $(\partial/\partial t)\alpha(r,t)$ is finite, observe that for any fixed t, $F(r,t)$ is sectionally continuous for $r \geq r_0$, so that the integral exists for $c \geq r_0$. Thus we may write

$$e^{2\alpha(r,t)} - 1 = e^{2\alpha(r,t) + \int_c^r F(r',t)dr'}, \quad (26)$$

whence

$$1 - e^{-2\alpha(r,t)} = e^{\int_c^r F(r',t)dr'}. \quad (27)$$

From (8) we have $b(r,t) = re^{\int_c^r F(r',t)dr'}$. So the requirement $b(r_0,t) = r_0$ now determines the arbitrary constant $c = r_0$. Thus,

$$b(r,t) = re^{\int_{r_0}^r F(r',t)dr'}. \quad (28)$$

Differentiating, we get for $r = r_0$

$$\frac{\partial}{\partial r}b(r_0,t) = e^{\int_{r_0}^{r_0} F(r,t)dr} + r_0 e^{\int_{r_0}^{r_0} F(r,t)dr}$$
$$\times \left[\frac{-2(\partial/\partial r)\gamma(r_0,t)}{K(r_0,t)(e^{2\alpha(r_0,t)} - 1)} - \frac{1}{r_0}\left(1 - \frac{1}{K(r_0,t)}\right)\right.$$
$$\left. - \frac{4e^{-\gamma(r_0,t)}e^{\alpha(r_0,t)}(\partial/\partial t)\alpha(r_0,t)}{e^{2\alpha(r_0,t)} - 1}\right]$$
$$= \frac{1}{K(r_0,t)} < 1, \quad (29)$$

by the assumption $K(r,t) > 1$.

The line element now becomes

$$ds^2 = -e^{2\gamma(r,t)}dt^2 + \frac{dr^2}{1 - e^{\int_{r_0}^r F(r',t)dr'}} + r^2\left(d\theta^2 + \sin^2\theta\, d\phi^2\right). \quad (30)$$

Since the flare-out conditions have been satisfied, the line element describes a traversable wormhole.

As an illustration, if α is time independent, $\gamma(r,t) \equiv 0$, and $K(r,t) = K$, a constant, then

$$e^{\int_{r_0}^r F(r',t)dr'} = e^{\int_{r_0}^r [-(1/r')(1-1/K)]dr'} = \left(\frac{r_0}{r}\right)^{1-1/K},$$
$$e^{2\alpha(r)} = \frac{1}{1 - (r_0/r)^{1-1/K}}, \quad (31)$$

which is Lobo's solution [20].

At this point, the following remark is in order: since we are only interested in the possible existence of wormholes, it is sufficient to note that to complete the description, the wormhole material should be cut off at some $r = a$ and joined to an external Schwarzschild spacetime. (See [20–22] for details.) This junction will make the space asymptotically flat, a critical feature referred to in the next section.

6. Implications

As we have seen, since $K(r,t) > 1$, line element (6) describes a traversable wormhole as long as the qualitative features in Figure 1 are met, resulting in a violation of the WEC. It is therefore conceivable that wormholes had formed spontaneously during the phantom-energy phase. Moreover, a possible mechanism for the formation of such wormholes is discussed in [23].

In this section, we study the consequences of our assumption that $K(r,t) \to 1$ sufficiently fast some time in the past. (As noted earlier, this limiting case corresponds to a cosmological constant.) Before doing so, however, let us recall that for an "arbitrary" wormhole the rate of change of the slope of γ is likely to be vanishingly small, leading to inequality (16). In addition,

$$\lim_{r\to r_0+}\alpha(r,t)=+\infty, \qquad \lim_{r\to r_0+}\frac{\partial}{\partial r}\alpha(r,t)=-\infty \quad \forall t, \quad (32)$$

while $\gamma=\gamma(r_0,t)$ must be finite to prevent an event horizon. Also, $\lim_{r\to\infty}\alpha(r,t)=\lim_{r\to\infty}\gamma(r,t)=0$ for all t. Finally, $(\partial/\partial r)\gamma(r,t)>0$ (see Figure 1).

In (22), as $K(r,t)\to 1$, we are left with

$$\frac{\partial}{\partial r}\alpha(r,t)+2e^{-\gamma(r,t)}e^{\alpha(r,t)}\frac{\partial}{\partial t}\alpha(r,t)=-\frac{\partial}{\partial r}\gamma(r,t). \quad (33)$$

As $r\to r_0+$, the left side goes to $-\infty$ (since $(\partial/\partial t)\alpha(r,t)$ is finite and nonpositive), so that the right side yields

$$\lim_{r\to r_0+}\frac{\partial}{\partial r}\gamma(r,t)=+\infty. \quad (34)$$

Since the gradient of the redshift function is related to the tidal forces [1], $(\partial/\partial r)\gamma(r_0,t)$ must remain finite. (In other words, γ must not have a vertical tangent at $(r_0,\gamma(r_0,t))$.) So for (34) to hold, $\gamma=\gamma(r,t)$ must approach the vertical line $r=r_0$ asymptotically. In other words, $\lim_{r\to r_0+}\gamma(r,t)=-\infty$, thereby producing an event horizon at the throat. The outcome is a black hole, since, in asymptotically flat spacetimes, a black hole is essentially characterized by the impossibility of escaping from a certain prescribed region to future null infinity. More formally, the boundary of $J^{-1}(\mathcal{I}^+)$, the causal past of future null infinity, must be a region that does not include the entire spacetime. But this requirement is met because our starting point is a well-defined region, the throat of a wormhole.

In summary, this paper discusses a wormhole geometry supported by a generalized form of phantom energy; the equation of state is $p(r,t)=-K(r,t)[\rho(r,t)+2g(r,t)]$, where $K(r,t)>1$ for all t. It is quite likely that the equation of state had crossed the phantom divide in the relatively recent past. If, in addition, the equation of state evolved so that $K(r,t)\to 1$ closer to the present, the result is a spacetime with a cosmological constant, a primary candidate for a model of the present Universe. It is shown that any existing wormhole would form an event horizon at the throat. Assuming that wormholes could have formed naturally during the phantom-energy phase, this outcome provides at least a *possible* explanation for the abundance of black holes and the complete lack of wormholes, even though wormholes are just as good a prediction of general relativity as black holes.

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