

Firewalls from General Covariance

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I define “horizon normalcy” as the approximate validity of semiclassical gravity and effective field theory for the description of observers that approach or cross a black hole horizon. If black holes return information, then horizon normalcy must fail substantially, at least in some global states. It has been proposed that horizon normalcy persists, so long as the Hawking radiation remains in a computationally simple state. Here I argue that state-dependent horizon normalcy—independent of the underlying mechanism and independent of the class of radiation states asserted to guarantee normalcy—requires a breakdown of general covariance far from the black hole, or else horizon normalcy will depend on the infinite future of the exterior. This is because the radiation can be in different states at different events, all spacelike to the horizon crossing event whose normalcy is at stake. I discuss a related effect in AdS/CFT, and I argue that its resolution by timefolds is of no help here.

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Firewall problem—The Almheiri-Marolf-Polchinski-Sully (AMPS) firewall paradox [1] can be formulated as follows [2]. A black hole is formed from a pure state and is allowed to evaporate fully. Assuming that (I) this process is described by a unitary S matrix, the entire Hawking radiation will be in some known pure state Ψ after evaporation is complete. (Evidence for unitarity includes [3–5]. The alternative, that information is lost [6], will not be explored here).

Consider an outgoing Hawking particle b in a neighborhood of a cut $v = v_0$ of the black hole horizon; see Fig. 1. For a Schwarzschild black hole of radius r_S , the near-horizon zone is the region $0 < r - r_S \ll r_S$. In this zone, the Hawking radiation modes are wave packets constructed from Rindler-like modes, with the black hole horizon playing the role of the Rindler horizon [7]. For definiteness, let b be a spherically symmetric wave packet, with radial support more than a Planck distance away from the horizon but with radial width much less than the black hole radius. We work in Eddington-Finkelstein coordinates,

$$ds^2 = -\left(1 - \frac{r_S}{r}\right)dv^2 - 2dv dr + r^2 d\Omega^2. \quad (1)$$

Now assume that (II) the evolution of b from the near-horizon region to the asymptotic region is governed by effective field theory. By assumptions (I) and (II), the

outgoing mode b will be entangled with and purified by the rest of the Hawking radiation in the state Ψ . (This is true in a typical state Ψ . If b is pure or less than typically entangled, the problem is worse [10]).

Not all of the Hawking radiation in Ψ has been emitted yet at the time $v = v_0$. However, at any time $v_D > v_P$, it is possible [11,12] to distill a nearly complete purification e_b of b from the “early” radiation R that has already come out by that time. Here $v_P \sim S_{\text{BH}} r_S$ is the Page time, when the coarse-grained Hawking radiation entropy begins to exceed the Bekenstein-Hawking entropy of the black hole. That is, there exists a (nonunique) unitary U acting on R at the time v_D that outputs a particle e_b such that $b e_b$ are together in a nearly pure state, tensored with some state for the remaining radiation system $R \setminus e_b$. The deviation from purity, the entropy $S_{b e_b}$, will be exponentially small in the number of Hawking particles emitted between v_P and v_D . We will take this number to be large, $T_H(v_D - v_P) \gg 1$, where $T_H \sim r_S^{-1}$ is the Hawking temperature.

If $v_D < v_0$, an observer Alice can put e_b in her pocket and schedule her jump so as to encounter b at $v = v_0$ near the horizon. By assumptions (I) and (II), Alice must find that b and e_b together are in the pure state predicted by unitarity. (Alice can verify this before she crosses the horizon, while she could still decide to stay outside and wait for the black hole to evaporate fully. This is important because black hole complementarity [13] allows for apparent contradictions between an infalling observer’s findings and the exterior description of a black hole, so long as the conflict cannot be experienced by or communicated to any single observer. Under this paradigm, if Alice were able to jointly probe b and e_b only *after* entering the black hole, she need not find them in the state demanded by unitarity).

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Let us now assume (III) horizon normalcy. Let N be a neighborhood of the event E of characteristic size r_S that excludes the high-curvature region near the singularity, but includes the near-horizon zone. Horizon normalcy is the statement that N is described by semiclassical gravity and effective field theory (SGEFT), up to nonperturbative corrections.

Horizon normalcy implies that the horizon itself looks locally like the vacuum [14]. In the vacuum state, b is strongly entangled with and purified by a “partner mode” $\tilde{b}$, which is also outgoing and localized just inside the horizon at the time v_0 [15]. (This step uses the fact that b is localized in the near-horizon region.)

To summarize, by (III), horizon normalcy, the von Neumann entropies satisfy $S_{\tilde{b}b} \ll 1$ and $S_b \sim O(1)$; by (I) and (II), unitarity, $S_{be_b} \ll 1$, but strong subadditivity [16] implies $S_{\tilde{b}b} + S_{be_b} \geq S_b$ [17]. Thus, Alice has local access to a tripartite system in a mathematically impossible quantum state. It follows that assumptions (I)–(III) cannot all be true [18].

If we commit to assumptions (I) and (II), then it follows that Alice will not see a normal horizon (III). This argument applies to any outgoing Hawking mode near the horizon, including higher angular momentum modes, limited only by some high-energy cutoff. Thus, the infalling observer should encounter a violent “firewall” [1].

Our focus will be on horizon normalcy, not on the violence of what might replace it. From a principled standpoint, any easily detectable departure from normalcy is as radical as a firewall. This includes the presence of just a single near-horizon mode $\tilde{b}b$ whose occupation number differs from the Unruh state, or the “nonviolent non-locality” models [21,22]. I will examine whether horizon normalcy can be recovered for a class of global states.

State-dependent horizon normalcy—The AMPS thought experiment is compelling because it is operational. When Alice crosses the horizon, she has local, simultaneous access to all three relevant particles $\tilde{b}, b, e_b$. Assuming (I) and (II), the fact that the horizon is abnormal (i.e., that at least $\tilde{b}b$ differs substantially from the vacuum state) is then beyond dispute.

In standard quantum mechanics, this conclusion holds regardless of whether anyone actually extracts e_b before Alice jumps in. Even if e_b remains scrambled in the radiation, the fact that it nearly purifies b makes it inconsistent for b to be purified by $\tilde{b}$ in the vacuum state.

Unless the purification e_b is decoded into a small system, however, this contradiction will not manifest itself directly to any one observer. (Alice cannot transport the entire early Hawking radiation to the horizon without creating a much larger black hole.) Moreover, the extraction of e_b is computationally hard: it requires a number of simple logical gates that scales exponentially with the area of the remaining black hole [20].

Based on these observations, it has been proposed that the horizon will remain normal if the state of the Hawking radiation remains in a certain class of states—for example, so long as no complex operation is performed on the Hawking radiation. This proposal implies nonlocality on the distance scale $r_S S_{\text{BH}}$ accessed by the Hawking radiation, which is much greater than that of the black hole. In addition, it requires—and so far, lacks—a precise definition of the normalcy-guaranteeing class of states. (For example, what counts as a “simple” operation? The rigorous definition of complexity classes takes the input size of the problem to infinity. Real-world settings are incompatible with this scaling, so a complexity class would have to be assigned at finite input size.) Moreover, one would need to keep track of the complexity of the “Hawking radiation” even if the latter is absorbed by other systems that could already be arbitrarily complex themselves.

We ignore these challenges here, in order to exhibit a robust, universal problem: state-dependent horizon normalcy requires a preferred global time slice and thus a violation of general covariance far from the black hole. Without a preferred time, horizon normalcy will depend on the infinite future—a form of nonlocality arguably too extreme to contemplate.

Time ambiguity—We begin by formalizing the idea of state-dependent horizon normalcy. Let Σ be a global Cauchy slice that contains E , and let ρ_N be the reduced state of the quantum fields in the open neighborhood N of E on Σ in which horizon normalcy should hold, extending less than $S_{\text{BH}} r_S$ from the horizon. Let $\rho_{\text{rad}}^{\text{enc}}$ (“encoded”) be the state of the radiation on Σ reached by the formation and natural evaporation of a black hole, at a distance of order $r_S S$ from the horizon. Let $\rho_{\text{rad}}^{\text{dec}}$ (“decoded”) be a radiation state on Σ in which e_b has been decoded. To obtain the state $\rho_{\text{rad}}^{\text{dec}}$ on Σ , some agent at an earlier time will have distilled the purification of b from the early radiation, so that e_b is available on Σ as a physical localized particle.

If the radiation state on Σ is $\rho_{\text{rad}}^{\text{dec}}$, then the AMPS argument precludes horizon normalcy: ρ_N must contain an excitation (at least) in the infalling mode built from b and $\tilde{b}$. But if the radiation state on Σ is $\rho_{\text{rad}}^{\text{enc}}$, then the above proposal asserts normalcy, i.e., it asserts that ρ_N is the vacuum state reduced to a neighborhood of E .

Now let Σ' be another Cauchy slice that coincides with Σ on N (and which therefore also contains E), but which meets the Hawking radiation at a different time; see Fig. 1. Thus, the (Schrödinger) state on Σ' can be different from that on Σ . For example, the radiation on Σ might be in the state $\rho_{\text{rad}}^{\text{enc}}$; but later e_b is decoded, so that on Σ' the state is $\rho_{\text{rad}}^{\text{dec}}$. Or conversely, e_b could have been decoded on Σ but then reencoded into the radiation by acting with $U^\dagger$, so that Σ' contains the state $\rho_{\text{rad}}^{\text{enc}}$.

If horizon normalcy near E depends on the state of the radiation, then it is not clear whether it should be determined from the radiation state on Σ , or on Σ' , or perhaps on

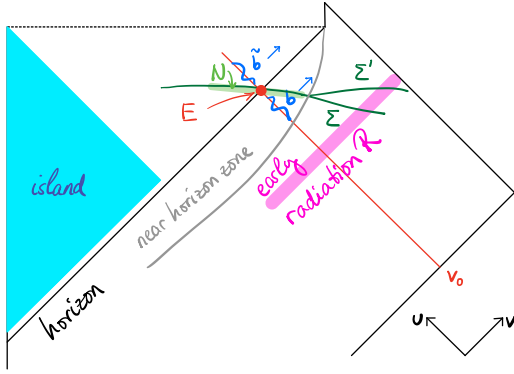


FIG. 1. Penrose diagram of an evaporating black hole. E is a fixed event at which Alice plans to cross the horizon. Here we examine the hypothesis [8,9] that horizon normalcy (the unambiguous validity of SGEFT in a local neighborhood N of E probed by Alice) depends on the state of distant Hawking radiation that has left the near-horizon zone before $v = v_0$. Since the radiation state can differ at different times Σ , Σ' , this hypothesis requires either the introduction of a preferred global time slice, or else the dependence of horizon normalcy at E on the infinite future of the exterior. Note that E lies outside of the entanglement island of the radiation; this is relevant for the discussion of timefolds.

yet another time slice. The only covariant choice is the boundary of the past of E , $v = v_0$ in our example. However, this choice fails [2]: Suppose the state at $v = v_0$ is $\rho_{\text{rad}}^{\text{enc}}$, leading to horizon normalcy at E . Then if e_b is decoded within a scrambling time $T_H^{-1} \log S_{\text{BH}}$ after v_0 , it can be sent into the black hole and received by Alice, leading again to an impossible quantum state of the tripartite system $\tilde{b}be_b$ locally accessible to Alice.

Therefore, horizon normalcy cannot depend on the state of the radiation, unless a preferred time slice is introduced far from the black hole on which that state is to be considered. However, this would violate general covariance far from the black hole, in a distant and weakly gravitating region. This seems far more radical than a breakdown of SGEFT only near the black hole, on a distance scale comparable to the black hole radius.

The remaining alternative is that horizon normalcy at E is guaranteed only if the radiation state remains in the horizon-normalcy-guaranteeing class (here, encoded) forever [23]. A dependence of normalcy at E on the infinite future of the exterior should again be contrasted with the comparatively mild nonlocality of a firewall.

The Maldacena-Susskind proposal and nonisometric maps—The time ambiguity problem exhibited above arises in any proposal in which horizon normalcy depends on the state of the distant Hawking radiation. It is irrelevant *why* horizon normalcy depends on the radiation state. It also does not matter whether the class of states sufficient for normalcy is characterized by the absence of complexity or by some other criterion. Nevertheless, it is instructive to

pinpoint how time ambiguity shows up in approaches that invoke a specific mechanism.

I will focus on two ideas that have received significant attention: the Maldacena-Susskind proposal [8,24] and nonisometric maps [9]. Both assert horizon normalcy if the radiation state lies in some restricted class, and thus both suffer from the time ambiguity. This is obscured, however, when we rely on toy models that capture key elements of each idea, but which omit too much of the physical spacetime structure of a real black hole to capture the time ambiguity.

The Maldacena-Susskind proposal invokes an analogy with a two-sided black hole in the Hartle-Hawking (or thermofield-double) state, which is assumed to have a smooth horizon. (This assumption is not trivial [31].) Assuming a typical entanglement pattern between the physical black hole and the radiation it has emitted by the time v_0 , there exists a unitary acting on the radiation that will convert the radiation into a second black hole of equal size, plus leftover radiation, with the two black holes in the thermofield-double state. Essentially, the Maldacena-Susskind proposal is the conjecture that this operation happens automatically as far as the infalling observer Alice is concerned, guaranteeing horizon normalcy. An exception has to be made if e_b has been decoded from the radiation, because then the horizon cannot be normal. This case is claimed [8] to be analogous to modifying the Hartle-Hawking state by removing or altering the mode $\tilde{b}$ when it is a simple physical wave packet near the left (unphysical) boundary of the two-sided black hole.

However, the decoding of e_b from the early Hawking radiation R is never simple. As a complex operation, it can be done, and undone, at various events that are all spacelike related to the horizon crossing event E . No criterion is available analogous to assessing whether $\tilde{b}$ is in the correct state on the left boundary at the (preferred) time when its state could be altered by a simple operator.

Let us turn to nonisometric maps [9], which implement state-dependent horizon normalcy as follows. One considers the fundamental state Ψ of the radiation and black hole (analogous to the boundary state in AdS/CFT or the final state of the radiation arriving at $\mathcal{I}^+$ in asymptotically flat spacetime). One asks if Ψ is the image of a simple effective global state ψ of the semiclassical geometry, under a nonisometric map V . If the answer is yes, then ρ_N is determined by reducing ψ to N . This ensures horizon normalcy if the Hawking radiation was not subjected to complex operations. If the answer is no, then horizon normalcy may fail in N . In particular, the answer will be no if e_b has been decoded from the radiation: since decoding requires a complex unitary, Ψ is then not the image of a simple effective state.

As in the Maldacena-Susskind proposal, the global state that is determinative for normalcy is implicitly assumed to

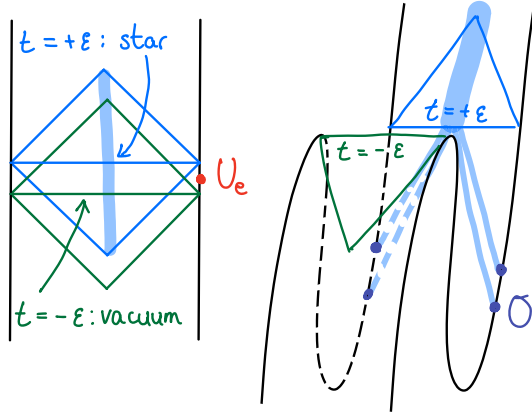


FIG. 2. Left: entanglement wedge conflict in AdS/CFT. At the boundary time $t = 0$, the boundary operator U_e changes the CFT vacuum to a state dual to a star. The Wheeler-DeWitt patches of the boundary times $t = \pm\epsilon$ (diamonds) overlap inconsistently in the bulk. Right: by inserting a timefold, the conflict is mitigated. The star can be built by sending in excitations from the boundary using local operators $\mathcal{O}$. There are two copies of the time $t = -\epsilon$: one in the vacuum state (green horizontal line) and one with a nearly complete star.

be unique and definite: the state is either simple or not. However, for a real black hole, we have seen that both statements can be true, on different Cauchy slices that all include the same neighborhood N of E . The key limitation of the toy model in Ref. [9] is that it fails to capture this aspect of causal structure (rather than, for example, that the map V was constructed only for a qudit model).

Entanglement wedge conflict and timefolds—Consider a holographic conformal field theory (CFT) in its vacuum state; see Fig. 2. At $t = 0$, act with a unitary U_e that maps the vacuum to an excited state, containing, for example, a star in the center of the bulk. U_e is nonlocal as a CFT operator. The entanglement wedge of any constant boundary time slice $t \neq 0$ is a well-defined Wheeler-DeWitt patch containing the entire bulk, either in the vacuum ($t < 0$) or with a star ($t > 0$). However, the entanglement wedges of $t = \pm\epsilon$ overlap for small ϵ and so are mutually inconsistent [32,33]. The sudden appearance of an excitation deep in the bulk, spacelike to U_e , would violate the Bianchi identity.

It is possible to mitigate entanglement wedge conflict by introducing a “timefold” [34]. The idea is to cut the boundary at $t = 0$ and “add in” a contour on which time runs backward and forward. This makes it possible to build up the new state at $t = \epsilon$ using local boundary operators $\mathcal{O}$ only. The length of the timefold and the operators $\mathcal{O}$ are highly nonunique.

The “time ambiguity” exhibited earlier has superficial commonalities with entanglement wedge conflict. The early Hawking radiation R would be analogous to the CFT, with ρ_{dec} and ρ_{enc} playing the role of the conflicting CFT states at $t = \pm\epsilon$. After the Page time, the entanglement

wedge of R contains an “island” that consists of the black hole interior a scrambling time earlier [4,5]. Hence Petz reconstruction of operators acting on $\tilde{b}$ succeeds [35,36], and changing the radiation state can change $\tilde{b}$. Might the conflicting black hole interiors dual to ρ_{dec} and ρ_{enc} be understood as different sheets on a timefold?

Yet, the analogy is imperfect: 1. State-dependent horizon normalcy [8,9] asserts that simple operators (in the sense of Ref. [20]) that act on *all* of the early radiation R will have no effect at the horizon, but certain complex operators will destroy horizon normalcy. However, the CFT operator U_e is simple (it arises by Hamiltonian evolution of local CFT operators over a polynomial time [37]) and yet has the large effect of creating the star. Indeed, almost every operation, simple or complex, on *all* of the star’s particles will have large effects. 2. No explicit construction of the timefold for Fig. 1, with E appearing on multiple sheets, is extant. The problem is subtle because Alice is never in the entanglement wedge of the radiation (the island); she enters only its future. (By contrast, any observers in the bulk of Fig. 2 will themselves be destroyed and recreated through the timefold.) The construction of a bulk timefold from two commuting boundary systems introduces additional ambiguities, which can be seen in the simpler setting of a “two-sided” Kruskal-AdS black hole. The left CFT plays the role of the radiation, and an observer enters the future of its entanglement wedge by falling into the black hole from the right. Decoding e_b corresponds in the two-sided setting to a boundary timefold on the left. A shock propagates from the left into the black hole on the second sheet; see

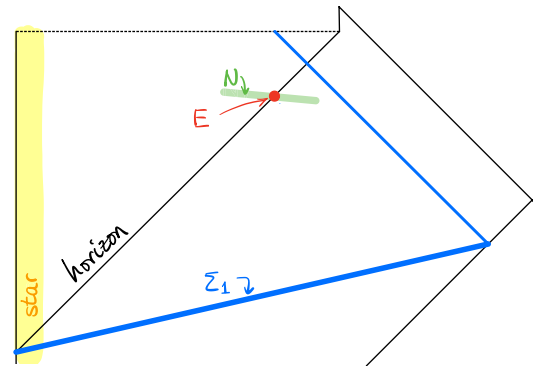


FIG. 3. Horizon normalcy requires the validity of SGEFT. Since N lies inside the domain of dependence of Σ_1 (blue wedge) according to SGEFT, the data on the partial Cauchy slice Σ_1 must evolve into the unique state in the region N predicted by SGEFT. This is inconsistent with timefolds, which assign different states to multiple copies of N . By contrast, the evolution from $t = -\epsilon$ to $t = +\epsilon$ in Fig. 2 is under control through the CFT and transcends SGEFT in the bulk. It can be given a bulk interpretation as a timefold; but there is no analog of the requirement of horizon normalcy, and hence no need for SGEFT to determine a unique bulk evolution between $t = \pm\epsilon$.

Eq. (4.4) and associated figures in Ref. [8]. However, the dual bulk timefold depends on the relative time ordering of the left and right CFT contours. This is unphysical because they are commuting operators. Moreover, if the left contour is executed first, then neither of the two sheets on the bulk timefold include the encounter of the infalling observer with the vacuum. This is because the entanglement wedge of the late-time left CFT state combined with the early-time right CFT state does not include a black hole interior at all.

We can put these unsolved issues aside, because timefolds would not restore horizon normalcy in any case. The question we set out to answer was: can horizon normalcy depend on the state of the radiation? Horizon normalcy requires the unambiguous validity of the standard predictions of SGEFT for the experiences of an observer crossing the horizon. Timefolds are not predicted [14], so horizon normalcy requires their absence; see Fig. 3. This, in turn, requires that no decoding operation will *ever* be performed on the radiation, again leading to a dependence of horizon normalcy on the infinite future.

Conclusion—For black holes to return information, new physics is needed at the horizon of a sufficiently old black hole. State-dependent horizon normalcy asserts that the new physics conspires to avoid detectability in some class of global states. I have shown that this necessitates a preferred time slicing far from the black hole; or else, horizon normalcy depends on the infinite future.

We have evidence for general covariance in weakly gravitating regions, but not for horizon normalcy, so it seems more likely that the new physics does not hide itself. Firewalls involve nonlocality “only” on the scale of the black hole. By conceding a universal failure of horizon normalcy, the firewall paradigm avoids the breakdown of SGEFT in weakly gravitating regions or even infinite nonlocality.

However, there is currently no proposal for an effective description of firewalls and their dynamics. They may be as far as $(Gr_s^2)^{1/4}$ from the horizon [38] in stationary black holes. A covariant criterion for their location will be suggested in forthcoming work.

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Data availability—No data were created or analyzed in this study.

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[18] This conclusion is independent of the underlying quantum gravity theory [19] from which SGEFT might emerge; it implies that (I)–(III) cannot all emerge. I omit a host of potential subtleties and suspected loopholes that were explored in the early literature on the firewall paradox (see Ref. [2] for a discussion and references). These include the hypothesis that the decoding unitary cannot be implemented fast enough, as a matter of physical law rather than mere practicality [20]. To my knowledge, this idea was not developed further; perhaps it should be.

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