

Excising Cauchy horizons with nonlinear electrodynamics

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Charged and/or rotating black holes in general relativity feature Cauchy horizons, which indicate a breakdown of predictability in the theory. Focusing on spherically symmetric charged black holes, we remark that the inevitability of Reissner-Nordström Cauchy horizon is due to the divergent electromagnetic self-energy of point charges. We demonstrate that any causal theory of nonlinear electrodynamics that regularizes the point charge self-energy also eliminates Cauchy horizons for weakly charged black holes. These black holes feature one (event) horizon and a spacelike singularity, analogous to the Schwarzschild metric. An example with Born-Infeld electrodynamics illustrates how this gives rise to an upper bound on the charge, which we compare to known limits.

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Introduction. Among all the effects that general relativity (GR) provides explanations for, *black holes* describe the most radical physical scenario, where an entire region becomes causally separated from the rest of the universe. The nature of causality in GR becomes ever more obscure and complex upon considering the maximal *analytical extension* of various black hole spacetimes. Major issues occur already when inspecting the extended spacetime causal structure in the simple case of adding an electric charge to the static spherically symmetric black hole, described by the Reissner-Nordström solution, or even the most astrophysically relevant case: the rotating black hole, described by the Kerr geometry.

Conjecturally—via cosmic censorship—GR has the grace of hiding pathologies it cannot further explain (curvature singularities, inner Cauchy horizons, and would-be chronology violations) behind horizons [1]. Cauchy horizons are especially pernicious as they mark a *loss of determinism*. Both Cauchy horizons and singularities signal a breakdown of the classical theory, yet the failure is worse at a Cauchy horizon. When evolution terminates at a spacelike singularity (as in Schwarzschild), the singular boundary does not contaminate the domain of

dependence in its past: the maximal Cauchy development is unique right up to that boundary. Moreover, curvatures at the singularity approach Planckian scales, going beyond the regime of validity of classical GR. Cauchy horizons, in contrast, can arise at sub-Planckian curvature, and the equations admit highly nonunique extensions compatible with identical past data. In this sense, singularities *localize* the breakdown of GR, whereas Cauchy horizons *globalize* it.

The problems with Cauchy horizons are addressed by Penrose’s *strong cosmic censorship* (SCC) conjecture [2]. In physical terms, SCC asserts that, generically, Cauchy horizons do not arise in physically reasonable solutions of GR. The inner horizons seen in, e.g., the Kerr-Newman family are then artifacts of fine-tuning; for generic data the causal structure should be Schwarzschild-like. The mass-inflation mechanism [3]—wherein perturbations are exponentially blueshifted at the Cauchy horizon—provides concrete support for SCC. This motivates the heuristic that GR “self-corrects”: perturbations drive mass inflation and replace the Cauchy horizon with a singular boundary, so classical evolution remains deterministic right up to where it fails.

SCC remains an open problem but what is understood points to a more subtle picture than what Penrose had envisioned [4,5]. However, irrespective of whether classical GR dynamically enforces SCC, an important physical question is whether the interior solutions—and their Cauchy horizons—should be trusted in the first place. Classical theories are expected to receive corrections of various kinds (e.g., quantum and higher derivative), and these are most pronounced in strong-field regions such as

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black hole interiors.¹ Such modifications could eliminate Cauchy horizons in parts of parameter space—thereby sidestepping the issues associated with SCC altogether—or introduce additional mechanisms that enforce it. Examples exist of the second possibility—in semiclassical gravity, the renormalized stress-energy generically diverges at the Cauchy horizon, triggering new instabilities [6]. These effects may be strong enough to restore SCC even in de Sitter settings, where the classical mechanism is on shakier ground. In this work we present examples of the first possibility: deformations of the classical theory that can eliminate the Cauchy horizon entirely.

We focus on charged, spherically symmetric black holes and ask how sensitive the Cauchy horizon is to *deformations* of the electromagnetic sector, while keeping Einstein gravity fixed. In linear Maxwell theory, the Coulomb self-energy of a point charge diverges, and we trace the inevitability of the Reissner-Nordström Cauchy horizon to this fact. Consequently, we consider theories of *non-linear electrodynamics* (NLE), motivated in part by their ability to *regularize* the self-energy of point sources [7,8]. Heuristically, one should expect NLE to play an important role in the causal structure of charged black holes: for weakly charged black holes, the field strength at the inner horizon can become very large, pushing the dynamics into the regime where NLE effects dominate. Our aim in this work is to uncover the implications of NLE in this regime.

Our main result is that, in any theory of NLE that satisfies the strong energy condition and has finite self-energy, the corresponding *weakly charged* black holes exhibit precisely the Schwarzschild causal structure—an event horizon and a spacelike curvature singularity. Since the standard energy conditions (null, weak, dominant, strong) are necessary for causal propagation in NLE, these results encompass the cases of physical interest [9]. More specifically, a Schwarzschild-like causal structure is realized whenever the gravitational mass exceeds the electromagnetic self-energy of the charge distribution. The weak-charge regime is also the most astrophysically relevant, because the net charge that black holes can accumulate is tightly constrained by both astrophysical and theoretical mechanisms [10–13]. Hence, for a broad class of physically viable models, the Cauchy horizon and its associated pathologies are avoided in practice.

To give a concrete example, we analyze the charged black hole solutions of Einstein-Born-Infeld theory [7], and show that the absence of the Cauchy horizon imposes an upper bound on the charge of a black hole. Given the known experimental and theoretical constraints on Born-Infeld theory, we discuss how this bound compares with theoretical and astrophysical limits.

¹Regular black hole models that remove the central singularity are often considered; we comment on them in Appendix A.

Single horizon charged black holes. NLE is a framework generalizing Maxwell’s classical electromagnetism by allowing the Lagrangian $\mathcal{L}$ to be a function of the two electromagnetic invariants,

$$\mathcal{L} = \mathcal{L}(\mathcal{S}, \mathcal{P}^2), \quad \mathcal{S} = \frac{1}{2} F_{\mu\nu} F^{\mu\nu}, \quad \mathcal{P} = \frac{1}{2} F_{\mu\nu} (*F^{\mu\nu}). \quad (2.1)$$

Here, $F_{\mu\nu} = (dA)_{\mu\nu}$ is the electromagnetic field strength, written in terms of the vector potential A_μ , and the Maxwell theory is recovered upon setting $\mathcal{L} = -\frac{1}{2}\mathcal{S}$.

When minimally coupled to GR, the corresponding generalized Einstein-Maxwell equations take the form,

$$\begin{aligned} dF &= 0, & d * D &= 0, & G_{\mu\nu} &= 8\pi T_{\mu\nu}, \\ D_{\mu\nu} &\equiv -2 \frac{\partial \mathcal{L}}{\partial F^{\mu\nu}} = -2(\mathcal{L}_\mathcal{S} F_{\mu\nu} + \mathcal{L}_\mathcal{P} (*F)_{\mu\nu}), \\ T^{\mu\nu} &= -\frac{1}{4\pi} (2F^{\mu\sigma} F^\nu{}_\sigma \mathcal{L}_\mathcal{S} + \mathcal{P} \mathcal{L}_\mathcal{P} g^{\mu\nu} - \mathcal{L} g^{\mu\nu}), \end{aligned} \quad (2.2)$$

where $\mathcal{L}_\mathcal{S} = \frac{\partial \mathcal{L}}{\partial \mathcal{S}}$ and $\mathcal{L}_\mathcal{P} = \frac{\partial \mathcal{L}}{\partial \mathcal{P}}$.

Consider a self-gravitating electrically charged spherically symmetric solution. In spherical symmetry, $\mathcal{P}$ identically vanishes and we have the restricted theory $\mathcal{L} = \mathcal{L}(\mathcal{S})$. Moreover, for any such theory the metric is necessarily described by a single function $f = f(r)$, e.g., [14], and can be written as

$$ds^2 = -f dt^2 + \frac{dr^2}{f} + r^2 d\Omega^2, \quad A = -\phi dt, \quad (2.3)$$

where $\phi = \phi(r)$, and $d\Omega^2 = d\theta^2 + \sin^2 \theta d\varphi^2$ is the area element on the sphere. The field equations then yield the following solution:

$$\begin{aligned} -2\mathcal{L}_\mathcal{S} E &= \frac{Q}{r^2}, & f &= 1 - \frac{2M}{r} + \frac{q^2(r)}{r^2}, \\ q^2(\rho) &\equiv -2\rho \int_\rho^\infty r^2 (2\mathcal{L}_\mathcal{S} E^2 + \mathcal{L}) dr, \end{aligned} \quad (2.4)$$

where $E = -\phi'(r)$. Here, Q is an integration constant related to the asymptotic electric charge by assuming that the field strength approaches that of Maxwell in the weak field limit (at large distances), i.e., $\mathcal{L}_\mathcal{S} \rightarrow -\frac{1}{2}$, giving $\lim_{r \rightarrow \infty} q(r) = Q$. Similarly, M is the “gravitational mass” integration constant. Since $\mathcal{S} = -E^2$, the first equation represents an algebraic equation for E .

An important observation in what follows is the connection between the function $q^2(r)$ appearing in the metric and the self-energy of a point particle in the specified theory of NLE. Explicitly, we define

$$U_{\text{self}}(r) = \frac{q^2(r)}{2r} = -\frac{1}{4\pi} \int T'_t dV, \quad (2.5)$$

which represents the electrostatic self-energy stored in a region between the radius r and infinity. The self-energy of a point source is given by $\lim_{r \rightarrow 0} U_{\text{self}}(r)$ [15,16].

In the case of Maxwell theory, the above evaluates to $U_{\text{self}} = Q^2/(2r)$. At short distances, the energy stored between r and infinity becomes large, blowing up in the $r \rightarrow 0$ limit. Hence, at sufficiently short distances the matter terms *necessarily* overwhelm the $2M/r$ gravitational potential, driving the metric function to $f(r \rightarrow 0) \rightarrow +\infty$. It is this short distance behavior of the self-energy that ultimately ensures that the Cauchy horizon is compulsory.

Let us now assume the following expansion of U_{self} near the origin $r = 0$:

$$U_{\text{self}} = \frac{U_{\text{self}}^{(-1)}}{r} + U_{\text{self}}^{(0)} + rU_{\text{self}}^{(1)} + O(r^2). \quad (2.6)$$

The absence of higher inverse powers of r indicates that we are assuming the NLE is not *more divergent* than Maxwell theory, for which the series contains only one nonzero term, $U_{\text{self}}^{(-1)} = Q^2/2$. In terms of the above we have at small r

$$f = \frac{2U_{\text{self}}^{(-1)}}{r^2} + \frac{2(U_{\text{self}}^{(0)} - M)}{r} + 1 + 2U_{\text{self}}^{(1)} + O(r). \quad (2.7)$$

For any Maxwell-like theory with $U_{\text{self}}^{(-1)} > 0$, Cauchy horizons are unavoidable. In such theories, adding Coulomb charge to Schwarzschild is nonperturbative: the change in the metric is not continuous.

Let us now assume finite self-energy, $U_{\text{self}}^{(-1)} = 0$ and $0 < U_{\text{self}}^{(0)} < \infty$, and study the black hole interior. The small- r behavior is controlled by the ratio of self-energy and mass, $U_{\text{self}}^{(0)}/M$. If $U_{\text{self}}^{(0)} < M$, then $f \rightarrow -\infty$ and the singularity is spacelike. If, in addition, the NLE satisfies the strong energy condition, there is generically *exactly one* horizon (see Appendix). This branch of solutions is therefore free of Cauchy horizons and we refer to it as the Schwarzschild-like (S) branch.

On the other hand, if $M < U_{\text{self}}^{(0)}$, the central singularity is timelike. If the NLE satisfies the strong energy condition, there is *at most* one inner horizon, so the causal structure is RN-like: either an event and inner horizon, or a naked timelike singularity. We call this the Reissner-Nordström-like (RN) branch.

Finally, the case $M = U_{\text{self}}^{(0)}$ is *marginal*: f remains finite at the origin, yet $r = 0$ is still a curvature singularity. Its character (spacelike, timelike, or null) is set by subleading terms in the expansion of U_{self} . Achieving a regular black hole further requires $U_{\text{self}}^{(1)} = 0$ so the core is Minkowski/(anti)-de Sitter (AdS), see however the no-go theorem in [17].

In all cases, $Q \rightarrow 0$ implies $U_{\text{self}}^{(0)} \rightarrow 0$, so the Schwarzschild metric is recovered. Weakly charged solutions therefore lie on the S-branch, with Schwarzschild-like causal and horizon structure—precisely because the self-energy remains finite at the origin.

An example: Born-Infeld electrodynamics. We now illustrate the previous section with the best-known NLE model, Born-Infeld theory [7]. It has many desirable features: finite point-charge self-energy, causality, no birefringence, electromagnetic duality, and a stress tensor satisfying the dominant and strong energy conditions [7,9,18–20]. Born-Infeld theory also arises as the low-energy effective action for gauge fields in string theory and D-brane physics [21,22].

The Lagrangian of Born-Infeld (BI) electrodynamics reads [7]

$$\mathcal{L}_{\text{BI}} = b^2 \left(1 - \sqrt{1 + \frac{\mathcal{S}}{b^2} - \frac{\mathcal{P}^2}{4b^4}} \right), \quad (3.1)$$

where b represents an upper bound on the field strength, with the Maxwell theory recovered in the limit $b \rightarrow \infty$. The corresponding self-gravitating solution reads

$$f = 1 - \frac{2M}{r} + \frac{2b^2}{r} \int_r^\infty \left(\sqrt{r^4 + \frac{Q^2}{b^2}} - r^2 \right) dr, \quad (3.2)$$

$$E = \frac{Q}{\sqrt{r^4 + Q^2/b^2}}.$$

The electromagnetic self-energy is finite and reads

$$U_{\text{self}} = b^2 \int_r^\infty \left(\sqrt{r^4 + \frac{Q^2}{b^2}} - r^2 \right) dr$$

$$= U_{\text{self}}^{(0)} - bQr + O(r^2),$$

$$U_{\text{self}}^{(0)} = \frac{1}{6} \sqrt{\frac{b}{\pi}} |Q|^{3/2} \Gamma\left(\frac{1}{4}\right)^2. \quad (3.3)$$

In the left panel of Fig. 1 we display the behavior of the metric function. We observe the S-branch (blue curve) with a single horizon and spacelike singularity for weakly charged black holes and the RN branch (red curve) with a timelike singularity and the Cauchy horizon (or naked singularity). The right panel of Fig. 1 illustrates the phase space of solutions as a function of the charge-to-mass ratio and the mass at fixed Born-Infeld parameter. The existence of S-type and RN-type black holes in Born-Infeld theory has long been known, e.g. [23,24], however the broader context and generality of this behavior described here is new.

Consider the RN-branch of solutions, which always terminates with an extremal black hole as the charge is

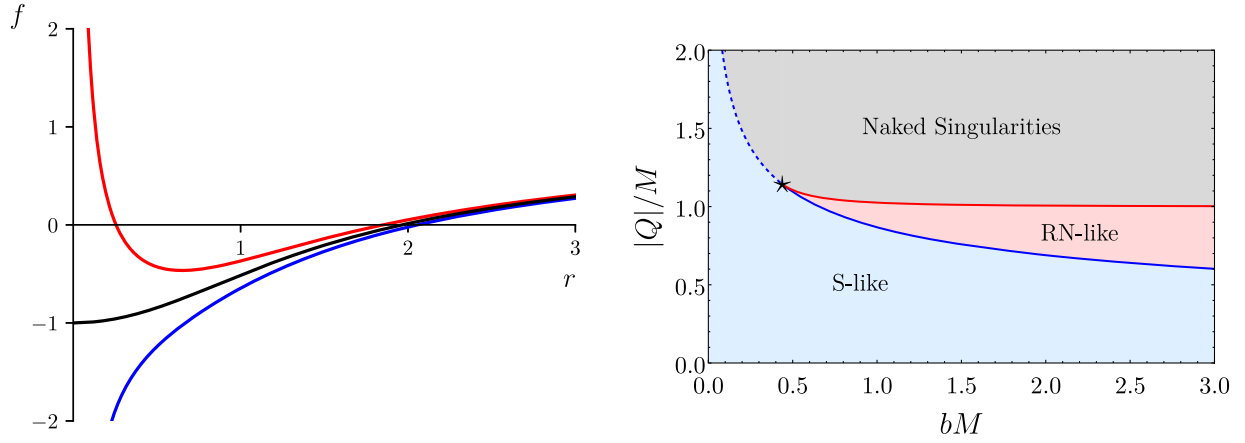


FIG. 1. Born-Infeld black holes. Left: Depending on (M, Q, b) there are three behaviors. The S branch ($M > U_{\text{self}}^{(0)}$; blue, lower) has a spacelike singularity and a single horizon. The RN branch ($M < U_{\text{self}}^{(0)}$; red, upper) has a timelike singularity and a Cauchy horizon. The marginal case ($M = U_{\text{self}}^{(0)}$; black, middle) separates the two. Right: Parameter space for the metric function f : S-like black holes (blue), RN-like black holes (red), and naked singularities (gray). The blue curve is the marginal line [equal to the bound (3.6)]; it is solid where marginal solutions are black holes and dashed where they are naked singularities. The solid red curve denotes extremal BI black holes. The star marks the point where marginal and extremal solutions coincide, yielding a null singularity.

increased. Extremal black holes in Born-Infeld theory have particular bounds on their charge-to-mass ratio that differ from Maxwell theory,

$$1 \leq \left(\frac{|Q|}{M} \right)_{\text{ext}} \leq \frac{6\sqrt{\pi}}{\Gamma(1/4)^2}. \quad (3.4)$$

The lower bound is achieved in the limit $bM \rightarrow \infty$ while the upper bound is exactly when the marginal solution is extremal, which requires

$$M = M_{\star} \equiv \frac{\Gamma(1/4)^2}{12\sqrt{2\pi}b}, \quad |Q| = Q_{\star} \equiv \frac{1}{2b}. \quad (3.5)$$

The RN branch ceases to exist for $M < M_{\star}$, for these smaller masses black holes are necessarily S type. In particular, note that S-type Born-Infeld black holes of *arbitrarily large* charge-to-mass ratio can exist, in contrast to the situation in Maxwell theory.

For $M > M_{\star}$, prohibiting the RN branch (or the Cauchy horizon) imposes an upper bound on the charge of a black hole derived from the requirement $M > U_{\text{self}}^{(0)}$,

$$\frac{|Q|}{M} < \left(\frac{6}{\Gamma(1/4)^2} \sqrt{\frac{\pi}{bM}} \right)^{2/3}. \quad (3.6)$$

On the other hand, for $M < M_{\star}$ this is precisely the bound required to ensure the absence of naked singularities. Hence for all masses, this provides a bound on the charge for which the black holes have physically sensible interiors. At fixed Born-Infeld parameter, this bound is tightest for large mass black holes while it is not particularly restrictive

in the limit of very small masses. For a fixed mass, the ratio tends to zero in the Maxwell limit $b \rightarrow \infty$, as required.

While cosmic censorship motivates the bound when $M < M_{\star}$, requiring it for $M > M_{\star}$ is in principle *ad hoc*. However, as we will now discuss, for physically viable values of the Born-Infeld parameter, the bound is actually very large by astrophysical standards. In their original paper, Born and Infeld estimated the value of b by requiring that the electron's self-energy is equal to its mass [7]. This estimate remains in line with the best modern constraints, e.g., [25], which require that $b \gtrsim 10^{20} \text{ V/m}^2$. Hence, converting to SI units we have

$$Q_{\text{max}} \approx 2.85 \left(\frac{M}{M_{\odot}} \right)^{2/3} \left(\frac{b_0}{b} \right)^{1/3} \times 10^{21} \text{ C}. \quad (3.7)$$

Here $M_{\odot}$ is one solar mass, while $b_0 = 10^{20} \text{ V/m}$ is taken as a reference scale for the Born-Infeld parameter.

A strict fundamental limit on charge accumulation arises from Schwinger pair production in the vicinity of the event horizon [11]. This will happen when the horizon electric field is $E \sim m_e^2/(\hbar e)$ in geometric units. This yields an upper bound on the charge $Q_{\text{max}}^{\text{pp}} \sim m_e^2 M^2/(\hbar e)$. If we equate pair production bound with the NLE bound, we find that $Q_{\text{max}}^{\text{NLE}} > Q_{\text{max}}^{\text{pp}}$ provided that

$$M \lesssim \left(\frac{\hbar^3 e^3}{m_e^6 b} \right)^{1/4} \sim 10^5 \left(\frac{b_0}{b} \right)^{1/4} M_{\odot}. \quad (3.8)$$

²In geometric units, $10^{20} \text{ V/m} \approx 9.59 \times 10^{-8} \text{ m}^{-1}$.

Hence, for all but the largest supermassive black holes, the inner horizon would be eliminated by NLE effects.

On the other hand, in any realistic environment the charge of black holes is well below the fundamental limit implied by the Schwinger effect. The bound is tightest for large black holes, so let us assess it for Sagittarius A*, which has $M \approx 4.3 \times 10^6 M_\odot$. Taking $b = b_0$, we find $Q_{\max}^{\text{Sgr A}^*} \approx 7.55 \times 10^{25}$ C. Comparing this with the observational constraints on the charge $Q_{\text{obs}}^{\text{Sgr A}^*} \lesssim 3 \times 10^8$ C and theoretical upper limit of $Q_{\text{theo}}^{\text{Sgr A}^*} \approx 10^{15}$ C [26], we see that both comfortably satisfy the bound. On the other hand, we may take the Planck scale as a firm upper limit to the value of b .³ If we take $b \sim 1/\ell_P$ where ℓ_P is the Planck length, then we find $Q_{\max}^{\text{Sgr A}^*} \approx 8.72 \times 10^{11}$ C, still well above the observational limit.

The main point of these considerations is not to make an experimental prediction, but to show the bound is physically reasonable. Avoiding Cauchy horizons requires that black holes are weakly charged, but this bound is still very large by astrophysical standards, making the Reissner-Nordström Cauchy horizon issue irrelevant in practice.

Although we illustrated these ideas with Born-Infeld electrodynamics, the conclusions are largely model independent. Any self-energy-regularizing NLE with strong-field scale b yields the same *parametric* scaling for the maximal charge as in (3.7). There is a simple physical reason for this: deviations from the Reissner-Nordström metric become important once the field strength at the inner horizon is comparable to b ,

$$\frac{Q}{r_-^2} \sim b. \quad (3.9)$$

Using $r_- = M - \sqrt{M^2 - Q^2}$ and expanding for large b , the Maxwell approximation is found to break down for

$$Q \lesssim M^{2/3} b^{-1/3}, \quad (3.10)$$

matching (3.7) up to order-one numerical factors. Moreover, if $Q \gg M^{2/3} \ell_P^{1/3}$, the curvature at r_- is sub-Planckian, so higher-curvature or quantum-gravity corrections should be negligible. Taking the Born-Infeld bound on b as representative, there is a wide window $9.59 \times 10^{-8} \text{ m}^{-1} \lesssim b \ll 6.19 \times 10^{34} \text{ m}^{-1}$ where the required separation of scales is realized for general NLEs.

Conclusions. Cauchy horizons are a pathological feature of Maxwell-charged black holes in general relativity. We showed that nonlinear electrodynamics with finite self-energy removes this pathology for weakly charged black

holes, yielding an effective upper bound on Q/M for physically viable solutions. Moreover, the conjecture of [27] suggests these S-branch black holes (unlike their RN-branch counterparts) are likely stable and admit a well-posed perturbation problem.⁴

The sensitivity of the Cauchy horizon to theory deformations has also been seen for asymptotically AdS black holes [28–31], e.g., in Einstein-Maxwell- Λ coupled to charged scalars or in deformations of Maxwell theory with additional massive gauge fields. Together with our results, this supports the view that the Reissner-Nordström Cauchy horizon is finely tuned, both under metric deformations *and* in theory space. Unlike those mechanisms, which rely on extra fields and AdS asymptotics, our conclusions arise in pure electrodynamics and are insensitive to whether the asymptotics are AdS, dS, or flat.

Our argument is not specific to electrodynamics, but to matter theories with finite self-energy. The main conclusions extend to any model with finite self-energy, a stress tensor satisfying $T^t_t = T^r_r$, and the strong energy condition. These are sufficient conditions, and the full situation is certainly more general than this. Motivated by this, one can consider more general geometries sourced by finite-self-energy matter, where the central singularity remains well behaved and spacelike whenever $M > M_m$, for some *marginal mass* M_m capturing the black hole’s nongravitational energy.

We have focused on charged, spherically symmetric spacetimes, and it is natural to ask how rotation changes the picture. In a companion paper [32], we studied Einstein-Born-Infeld theory and found that the Cauchy horizon is removed for three-dimensional rotating black holes and for four-dimensional rotating black-branes. We also presented strong evidence that nonlinear electrodynamics can eliminate the Kerr-Newman Cauchy horizon in some parameter regimes.

Could an analogous mechanism act in the astrophysically relevant Kerr spacetime, replacing its Cauchy horizon with a spacelike singularity? There is reason for optimism in realistic collapse: although the exterior rapidly approaches the Kerr metric, the interior can remain sensitive to the collapsing matter, and such matter effects could plausibly remove the Cauchy horizon. The vacuum case is harder. Since the Kerr Cauchy horizon is not tied to matter self-energy, eliminating it in pure gravity would likely require modified gravitational dynamics. Still, because the rotational contribution has a Maxwell-like structure,

$$\frac{\Delta}{r^2} = 1 - \frac{2(M - U)}{r}, \quad U = \frac{a^2}{2r}, \quad (4.1)$$

³In stringy Born-Infeld theory, $b \sim \ell_s^{-1}$, where ℓ_s is the string length. Although ℓ_s is often taken to be near the Planck length, this need not be the case.

⁴Interestingly, eliminating Cauchy horizons altogether—including in the strongly charged regime—is impossible within a pathology-free NLE framework.

it is tempting to speculate that a similar construction might work. Such metrics are easy to write down off shell (e.g., [33]), but it remains unclear whether they arise from any consistent modified (e.g., limiting-curvature) gravity theory.

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

Appendix A: Regular black holes. As noted in the introduction, both singularities and Cauchy horizons signal a breakdown of GR. One popular response is the regular black hole (RBH) program [34–36], where one postulates nonsingular metrics and studies their geometry. However, in the static, spherically symmetric sector, RBH metrics necessarily contain Cauchy horizons. This is a serious obstruction: since the metric admits no unique extension across a Cauchy horizon, the region “beyond” it is not well-defined physically, so one cannot robustly claim the singularity has been resolved.

The RBH program can remain consistent with predictivity so long as the inner horizon is not a Cauchy horizon. This necessitates going beyond static metrics and considering dynamical and evaporating ones [37–39]. However, in this case the instabilities that enforce SCC in GR take on a more problematic role. Proposed cures for mass inflation include making the inner horizon degenerate/extremal [40], but this appears nongeneric: it demands *ad hoc* parameter tuning, and extremal horizons suffer additional instabilities [41]. Moreover, while inner extremality may suffice to avoid classical mass inflation, it may still fail to evade semiclassical instabilities [42]. Ultimately, the fate of the inner horizon of RBH models is a dynamical question that will depend on the particular model in which RBHs are obtained as solutions. In any case, in a regime where classical dynamics is driven by GR, a Schwarzschild-like causal structure with spacelike curvature singularity is the more conservative possibility, avoiding the more pernicious pathologies.

Appendix B: Proof of horizon structure. If an NLE violates any energy condition, the theory is acausal [9]. We therefore study how these inequalities constrain the interior structure, focusing on the strong energy condition (SEC), which is satisfied by, e.g., Born-Infeld and RegMax electrodynamics. In the main text we showed that finite self-energy NLEs admit only Schwarzschild-like or RN-like small- r asymptotics. One might nevertheless try to engineer an NLE with these asymptotics but with arbitrarily many additional inner horizons, see, e.g., [43]. We will show this is impossible for causal NLEs: extra inner horizons generically force a violation of the SEC.

Using (2.4) and (2.5), the metric function can be written in terms of the self-energy,

$$f(r) = 1 - \frac{2M}{r} + \frac{2U(r)}{r}. \quad (\text{B1})$$

The SEC requires $(T_{\mu\nu} - \frac{1}{2}Tg_{\mu\nu})u^\mu u^\nu \geq 0$ for all timelike u^μ . Restricting to radial timelike curves and using the field equations gives

$$\text{SEC} \Rightarrow U''(r) \geq 0. \quad (\text{B2})$$

Assume finite self-energy with $U_{\text{self}}^{(0)} < M$, so the interior is Schwarzschild-like: $f(r \rightarrow 0) \rightarrow -\infty$ while $f(r \rightarrow \infty) \rightarrow 1$. Define

$$g(r) \equiv rf(r) = r - 2M + 2U(r), \quad (\text{B3})$$

so f and g have the same zeros. Since $g''(r) = 2U''(r)$, the SEC implies $g''(r) \geq 0$, i.e. g is convex. Moreover $g(r \rightarrow 0) \rightarrow U_{\text{self}}^{(0)} - M < 0$ and $g(r \rightarrow \infty) \rightarrow \infty$. A convex function with these endpoint behaviors can cross zero at most once, so no inner horizon can exist.

For RN-like small- r behavior, $U_{\text{self}}^{(0)} > M$ and hence $g(r \rightarrow 0) \rightarrow U_{\text{self}}^{(0)} - M > 0$. Convexity then allows at most two zeros (one may be degenerate), corresponding to an event and a single inner horizon. Therefore, any additional horizons beyond one in the S branch or two in the RN branch generically violate the SEC, and hence causality in the NLE. The physically viable interior causal structures are thus those described in the main text.

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