The quantum physics of chronology protection

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Abstract:
This is a brief survey of the current status of Stephen Hawking’s “chronology protection conjecture”. That is: “Why does nature abhor a time machine?” I’ll discuss a few examples of spacetimes containing “time machines” (closed causal curves), the sorts of peculiarities that arise, and the reactions of the physics community. While pointing out other possibilities, this article concentrates on the possibility of “chronology protection”. As Stephen puts it:

It seems that there is a Chronology Protection Agency which prevents the appearance of closed timelike curves and so makes the universe safe for historians.

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Simply put, chronology protection is the assertion that nature abhors a time machine. In the words of Stephen Hawking [1]:

*It seems that there is a Chronology Protection Agency which prevents the appearance of closed timelike curves and so makes the universe safe for historians.*

The idea of chronology protection gained considerable currency during the 1990’s when it became clear that traversable wormholes, which are not too objectionable in their own right [2, 3, 4], seem almost generically to lead to the creation of time machines [5, 6, 7, 8]. The key word here is “seem”. There are by now many technical discussions available in the literature (well over 200 articles), and in the present chapter I will simply give a pedagogical and discursive overview, while adding an extensive bibliography for those interested in the technical details. First: a matter of language, for all practical purposes the phrases “time machine” and “closed timelike curve” (or the closely related “closed null curve”) can be used interchangeably.

**Why is chronology protection even an issue?**

Before embarking on a discussion of chronology and how it is believed to be protected [1, 6, 7], it is useful to first ask why chronology even needs to be protected. In Newtonian physics, and even in special relativity or flat-space quantum field theory, notions of chronology and causality are so fundamental that they are simply built into the theory ab initio. Violation of normal chronology (for instance, an effect preceding its cause) is so objectionable an occurrence that any such theory would immediately be rejected as unphysical.
Unfortunately, in general relativity one cannot simply assert that chronology is preserved, and causality respected, without doing considerable additional work. The essence of the problem lies in the fact that the Einstein equations of general relativity are local equations, relating some aspects of the spacetime curvature at a point to the presence of stress-energy at that point. Additionally, one also has local chronology protection, inherited from the fact that the spacetime is locally Minkowski (the Einstein Equivalence Principle), and so “in the small” general relativity respects all of the causality constraints of special relativity.

What general relativity does not do is to provide any natural way of imposing global constraints on the spacetime — certainly the Einstein equations provide no such nonlocal constraint. In cosmology this leads to the observation that the global topology of space is not constrained by the Einstein equations; spatial topology is an independent discrete variable that has to be decided by observation. (And this requires additional data over and above whatever is needed to decide the familiar $k = +1$, $k = 0$, or $k = -1$ question of the Friedmann–Robertson–Walker cosmologies [9].) Similarly, global temporal topology is not constrained by the Einstein equations themselves, and additional physical principles need to be brought into play to somehow deal with the possibility of nontrivial-temporal topology.

Without imposing additional principles along these lines general relativity is completely infested with time machines (in the sense of closed causal curves). Perhaps the earliest examples of this pathology are the Van Stockum spacetimes [10], but the example that has attracted considerably more attention is Kurt Gödel’s peculiar cosmological solution [11]. These spacetimes are exact solutions of the Einstein equations, with sources that (at least locally) look physically reasonable, which nevertheless possess serious global pathologies. If it were only a matter of dealing with these two particular examples, physicists would not be too worried — but similar behaviour occurs in many other geometries, for instance, deep inside the Kerr solution. A complete list of standard but temporally ill-behaved spacetimes is tedious to assemble, but at a minimum should include:

1. Gödel’s cosmology [11];
2. Van Stockum spacetimes [10]/ Tipler cylinders [12]/ longitudinally spinning cosmic strings [8];
4. Gott’s time machines [14];
5. Wheeler wormholes (spacetime foam) [15, 16, 17];
6. Morris–Thorne traversable wormholes [2, 5];

7. Alcubierre “warp drive” spacetimes [18].

The Wheeler wormholes are included based on theorems that localized topology change implies either causal pathology or naked singularities; either possibility is objectionable [19, 20, 21]. The Morris–Thorne traversable wormholes are included based on the observation that apparently trivial manipulations of these otherwise not too objectional geometries seem to almost generically lead to the development of closed timelike curves and the destruction of normal chronology [5, 8]. For the “warp drive” spacetimes manipulations similar to those performed for traversable wormhole spacetimes seem to lead inevitably to time travel. (Once one has effective faster-than-light travel, whether via wormholes or warpdrives, the twin pseudo-paradox of special relativity is converted into a true paradox, in the sense of engendering various time travel paradoxes.)

Now in each of these particular cases you can at a pinch find *some* excuse for not being too concerned, but it’s a different excuse in each case. The matter sources for the Gödel solution are quite reasonable, but the observed universe simply does not have those features. The Van Stockum time machines and their brethren require infinitely long cylindrical assemblages of matter rotating at improbable rates. The Kerr and Kerr–Newman pathologies are safely hidden behind the Cauchy horizon [13], where one should not trust naive notions of maximal analytic extension. (The inner event horizon is classically unstable.) The Wheeler wormholes (spacetime foam) have never been detected, and at least some authors now argue against the very existence of spacetime foam. The energy condition violations implicit in traversable wormholes and warp drive spacetimes do not seem to be qualitatively insurmountable problems [2, 8], but do certainly give one pause [22]. This multiplicity of different excuses does rather make one worry just a little that something deeper is going on; and that there is a more general underlying theme to these issues of (global) chronology protection.

**Paradoxes and responses.**

Most physicists view time travel as being problematic, if not downright repugnant. There are two broad classes of paradox generated by the possibility of time travel, either one of which is disturbing:

1. Grandfather paradoxes: Caused by attempts to “change the past”, and so modify the conditions that lead to the very existence of the entity that is trying to “modify the timestream”.

2. Bootstrap paradoxes: Where an effect is its own cause.
Faced with the *a priori* plethora of geometries containing closed timelike curves, with the risk of these two classes of logical paradox arising, the physics community has developed at least four distinct reactions [8]:

1. Make radical alterations to our worldview to incorporate at least some versions of chronology violation and “time travel”. (The “radical re-write” conjecture.) One version of the radical re-write conjecture uses non-Hausdorff manifolds to describe “train track” geometries where the same present has two or more futures (or two or more pasts). A slightly different version uses the “many worlds” interpretation of quantum mechanics to effectively permit switching from one history to another [23]. More radically one can even contemplate multiple coexisting versions of the “present”.

2. Permit constrained versions of closed timelike curves — supplemented with a consistency condition that essentially prevents any alteration of the past. (This is the essence of the Novikov consistency conditions [24, 25, 26].) The consistency conditions are sometimes summarized as “you can’t change recorded history” [27]. The central idea is that there is a single unique timeline so that even in the presence of closed timelike curves there are constraints on the possibilities that can occur. In idealized circumstances these consistency constraints can be derived from a least action principle. More complicated situations seem to run afoul of the notion of “free will”, though there is considerable doubt as to the meaning of “free will” in the presence of time travel [28].

3. Appeal to quantum physics to intervene and provide a universal mechanism for preventing the occurrence of closed timelike curves. This, in a nutshell, is Stephen Hawking’s “chronology protection” option, the central theme of this chapter, which we shall develop in considerable detail below.

4. Agree to not think about these issues until the experimental evidence becomes overwhelming. (The “boring physics” conjecture.) After all, what is the current experimental evidence? Assume global hyperbolicity and cosmic censorship and be done with it. If, for instance, one takes canonical gravity seriously as a fundamental theory then there exists at least one universal foliation by complete spacelike hypersurfaces. This automatically forbids closed timelike curves at the kinematical level, before dynamics (classical or quantum) comes into play. However, it should be noted that canonical gravity interpreted in this strict sense has severe difficulties (for instance, in dealing with maximal analytic extensions of the Kerr spacetime).
Originally it was hoped that it would be possible to decide between these options based on classical or at worst semiclassical physics — however it is now becoming increasingly clear that the ultimate resolution of the chronology protection issue will involve deep issues of principle at the very foundations of the full theory of quantum gravity.

**Elements of chronology protection.**

Chronology protection is at one level an attempt at “having one’s cake and eating it too” — this in the sense that it provides a framework sufficiently general to permit interesting and nontrivial topologies and geometries, but seeks to keep the unpleasant side effects under control. Chronology protection deals with the localized production and destruction of closed timelike curves; the very essence of what we might like to think of as “creating” a time machine.

(Cosmological time machines, in the sense of Gödel, are best viewed as an example of the GIGO principle; garbage in, garbage out. Just because one has a formal solution to a set of differential equations does not mean there is any physical validity to the resulting spacetime. A differential equation without boundary conditions/initial conditions has little predictive power, and it is very easy to generate ill-posed problems. Cosmological time machines are by definition intrinsically and equally sick everywhere in the spacetime.)

In the case of a localized production of closed timelike curves the situation is more promising: the spacetime is then divided into regions of normal causal behaviour and abnormal causal behaviour, with the boundary that separates these regions referred to as the “chronology horizon”. It is the behaviour of quantum physics at and near this chronology horizon that provides the basis for chronology protection.

Specifically, a point \( x \) is part of the chronology violating region if there is a closed causal curve (closed timelike curve) or closed chronological curve (closed null/timelike curve) passing through \( x \). The chronology horizon is then defined as the boundary of the future of the chronology violating region. (That is, the boundary of the region from which chronology violating physics is visible.) This chronology horizon is by definition a special type of Cauchy horizon. Under reasonably mild technical conditions Hawking has argued that the chronology horizons appropriate to locally constructed time machines should be compactly generated and contain a “fountain”; essentially the first closed null curve to come into existence as the time machine is formed [1].

A classical photon placed on this fountain will circulate around the fountain infinitely many times; in effectively zero “elapsed” time. On each circuit around the fountain there is generically a nontrivial holonomy
that changes the energy of the photon. For a past chronology horizon, which expands as we move to the future (as defined by someone outside the chronology violating region) this provides a boost, a net increase in the photon energy for each circuit of the fountain. The photon energy increases geometrically, reaching infinity in effectively zero time [1]. On each circuit

\[ E \rightarrow e^h E \rightarrow e^{2h} E \ldots; \quad h = -2 \oint \mathcal{R}(\epsilon) \, dt, \]

with the size of the energy boost being controlled by a loop integration around the fountain involving the Newman–Penrose parameter \( \epsilon \). (In simple situations involving wormholes this holonomy is essentially the Doppler shift factor due to relative motion of the wormhole mouths, but when phrased in terms of \( \oint \epsilon \) it can be generalized to arbitrary chronology horizons possessing a fountain.) The source of this energy must ultimately be the spacetime geometry responsible for the chronology horizon, and by extension, the stress-energy used to warp spacetime and set up the fountain in the first place. If we now let the photon (and the gravitational field it generates) back-react on the spacetime, its infinite energy will presumably alter the spacetime geometry beyond all recognition.

Unfortunately this is a classical argument, appropriate to a classical point particle following a precisely defined null curve. Will quantum physics amplify or ameliorate this effect? Real photons are wave-packets with a certain transverse size, and generically the same effect that leads to the energy being boosted leads to the wave-packet being defocussed — the geometry in a tubelike region surrounding the fountain acts as a diverging lens [1].

With two competing effects, the question becomes which one wins? The answer, “it depends”. There are geometries for which the classical defocussing effect overwhelms the boost effect, and the classical stress tensor remains finite on the fountain. There are other geometries for which the reverse holds true. But this certainly means that classical effects do not provide a universal mechanism for eliminating all forms of closed causal curves. Thus the search for a universal chronology protection mechanism must then (at the very least) move to the semiclassical quantum realm.

**Semiclassical arguments**

In semiclassical quantum gravity, one treats gravity as a classical external field, but one quantizes everything else. So far, this is just curved space quantum field theory. But then one additionally demands that the Einstein equations hold for the quantum expectation value of the stress-
energy tensor:
\[ G_{\mu\nu} = 8\pi G_{\text{Newton}} \langle \psi | T_{\mu\nu} | \psi \rangle. \]

Semiclassical quantum gravity seems [at first glance] to lead to a universally true statement to the effect that the renormalized expectation value of the stress-energy tensor blows up at the chronology horizon. The idea is based on the fact that in curved manifolds (modulo technical issues to be discussed below) the two-point correlation function (Green function; a measure of the mean square fluctuations) of any quantum field is of Hadamard form
\[
G(x,y) = \sum_{\gamma} \frac{\Delta_{\gamma}(x,y)^{1/2}}{4\pi^2} \left\{ \frac{1}{\sigma_{\gamma}(x,y)} + \nu_{\gamma}(x,y) \ln |\sigma_{\gamma}(x,y)| + \varpi_{\gamma}(x,y) \right\}.
\]

Here the sum runs over the distinct geodesics from \( x \) to \( y \); the quantity \( \Delta_{\gamma}(x,y) \) denotes the Van Vleck determinant evaluated along the geodesic \( \gamma \); the quantity \( \sigma_{\gamma}(x,y) \) denotes Synge’s “world function” (half the square of the geodesic distance from \( x \) to \( y \)); and the two functions \( \nu_{\gamma}(x,y) \) and \( \varpi_{\gamma}(x,y) \) are smooth with finite limits as \( y \to x \). Provided the Green function can be put into this Hadamard form, the expectation value of the point split stress-energy tensor can be defined by a construction of the type
\[
\langle T_{\mu\nu}(x,y,\gamma_0) \rangle = D_{\mu\nu}(x,y,\gamma_0) G(x,y).
\]
Here \( \gamma_0 \) denotes the trivial geodesic from \( x \) to \( y \) (which collapses to a point as \( y \to x \), this geodesic will be unique provided \( x \) and \( y \) are sufficiently close to each other), while \( D_{\mu\nu}(x,y,\gamma_0) \) is a rather complicated second-order differential operator built up out of covariant derivatives at \( x \) and \( y \). The covariant derivatives at \( y \) are parallel transported back to \( x \) along the geodesic \( \gamma_0 \) with the result that \( \langle T_{\mu\nu}(x,y,\gamma_0) \rangle \) is a tensor with respect to coordinate changes at \( x \), and a scalar with respect to coordinate changes at \( y \). One then defines the renormalized expectation value of the stress-energy tensor by taking the limit \( y \to x \) and discarding the universal divergent piece which arises from the contribution of the trivial geodesic to the Green function. In other words, the renormalized Green function is defined by
\[
G(x,y)_R = \sum_{\gamma \neq \gamma_0} \frac{\Delta_{\gamma}(x,y)^{1/2}}{4\pi^2} \left\{ \frac{1}{\sigma_{\gamma}(x,y)} + \nu_{\gamma}(x,y) \ln |\sigma_{\gamma}(x,y)| + \varpi_{\gamma}(x,y) \right\},
\]
and the renormalized stress energy by
\[
\langle T_{\mu\nu}(x) \rangle_R = \lim_{y\to x} D_{\mu\nu}(x,y,\gamma_0) G_R(x,y).
\]
Other methods of regularizing and renormalizing the stress-energy could be used, the results will qualitatively remain the same. The net result is that

\[
\langle T_{\mu\nu}(x) \rangle_R = \sum_{\gamma \neq \gamma_0} \frac{\Delta_\gamma(x, x)^{1/2}}{\sigma_\gamma(x, x)^2} \cdot t_{\mu\nu}(x) + \cdots
\]

Here \( t_{\mu\nu}(x) \) is a dimensionless tensor built up out of the metric and tangent vectors to the geodesic \( \gamma \), while the \( \cdots \) denote subdominant contributions. The key observation is that if any of the non-trivial geodesics from \( x \) to itself are null (invariant length zero), then there is an additional infinity in the stress-energy over and above the universal local contribution that was removed by renormalization. (For a slightly different way of doing things, one could just as easily choose to work with the effective action \[29\] instead of the stress-energy; the conclusions are qualitatively similar.)

In general these self-intersecting null geodesics define the \( N \)'th-polarized hypersurfaces, where \( N \) is a winding number which counts the number of times the geodesic passes through the tubular region surrounding the fountain. These polarized hypersurfaces lie inside the chronology horizon and typically approach it as \( N \to \infty \,[6, 7] \). In particular, the fountain is a nontrivial closed null geodesic, and this argument indicates that the renormalized stress-energy tensor diverges at the fountain. But infinite stress-energy implies, via the Einstein equations, infinite curvature. The standard interpretation of this is (or rather, was) that once back-reaction is taken into account the fountain (and \textit{ipso facto}, the entire chronology horizon) is destroyed by the (mean square) quantum fluctuations. (You do not need the stress-energy to diverge everywhere on the chronology horizon; it is sufficient if it diverges at the fountain.)

The fly in the ointment here is these same quantum fluctuations. On the one hand the quantum fluctuations are responsible for the formal infinity in the expectation value of the stress-energy at the fountain, on the other hand: Does the back-reaction due to the expectation value of the stress-energy tensor become large before the quantum fluctuations in the metric completely invalidate the manifold picture? (This very question led to a spirited debate between Stephen Hawking and Kip Thorne \[6, 7\], with disagreement on how to define the notion of “closeness” to the chronology horizon.)

It is now generally accepted that typically the back reaction becomes large before metric fluctuations invalidate the manifold picture, but that there are exceptional geometries where the back-reaction can be kept arbitrarily small arbitrarily close to the chronology horizon. A particular example of this phenomenon is if you take a “ring configuration” of wormholes, where each individual wormhole is nowhere near forming a
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chronology horizon, but the combination is just on the verge of violating causality [30]. Then there is a closed spacelike geodesic which traverses the entire ring of wormholes whose invariant length is becoming arbitrarily small; but because the spacelike geodesic is traversing many wormhole mouths (each of which acts as a defocussing lens) the Van Vleck determinant can be made arbitrarily small in compensation.

That is: adopt the length of the shortest closed spacelike geodesic as a diagnostic for how close the spacetime is to forming a time machine. Then no matter how close one is to violating chronology, there are some geometries for which the renormalized stress-energy tensor (and the quantum-induced back reaction) can be made arbitrarily small. In a similar vein there are a number of other special case examples (for example, toy models based on variants of the Grant and Misner spacetimes [31, 32, 33]) for which the renormalized stress-energy remains finite all the way up to the chronology horizon. The upshot of all this is that the search for a universal chronology protection mechanism must (at the very least) involve issues deeper and more fundamental than the size of the quantum-induced back reaction.

The failure of semiclassical gravity

The most mathematically precise and general statements known concerning the nature of the pathology encountered at the chronology horizon are encoded in the singularity theorems of Kay, Radzikowski, and Wald [34]. In a highly technical article using micro-local analysis they demonstrated:

**Theorem 1** There are points on the chronology horizon where the two-point function is not of Hadamard form.

Because there are points where the two-point function is not of Hadamard form, the entire process of defining a renormalized stress-energy tensor breaks down at those points. That is:

**Corollary 1** There are points on the chronology horizon where semiclassical Einstein equations fail to hold.

Note that the semiclassical Einstein equations,

\[ G_{\mu\nu} = 8\pi G_{\text{Newton}} \langle T_{\mu\nu} \rangle_R, \]

fail for a subtle reason; they fail simply because at some points the RHS fails to exist, not necessarily because the RHS is infinite. Now typically, based on the explicit calculations of the last section, the renormalized stress-energy does blow up on parts of the chronology horizon. The significant new feature of the Kay–Radzikowski–Wald analysis is that even if
the stress-energy remains finite as one approaches the chronology horizon, there will be points on the chronology horizon for which no meaningful limit exists. (For a specific example, see [35].)

The physical interpretation is that semiclassical quantum gravity fails to hold (at some points) on the chronology horizon; a fact which can be read in two possible ways:

1. If you assume that semiclassical quantum gravity is the fundamental theory (at best a minority opinion, and there are good very reasons for believing that this is not the case), then by *reductio ad absurdum* the chronology horizon must fail to form. Chronology is protected, essentially by *fiat*.

2. If you are willing to entertain the possibility that semiclassical quantum gravity is not the whole story (the majority opinion), then it follows from the above that issues of chronology protection cannot be settled at the semiclassical level. Chronology protection must then be settled (one way or another) at the level of a full theory of quantum gravity.

An attractive physical picture that captures the essence of the situation is this: Sufficiently close to (but outside) the chronology violating region there are extremely short self-intersecting spacelike geodesics. The length of these geodesics can be used to develop an observer independent measure of closeness to chronology violation. Indeed let

\[ \mathcal{M}(\ell) = \left\{ x \mid \exists \gamma \neq \gamma_0 : \sigma_\gamma(x, x) \leq \frac{\ell^2}{2} \right\} \]

Then \( \mathcal{M}(0) \) is one way of characterizing the chronology violating region, while \( \mathcal{M}(L_{\text{Planck}}) - \mathcal{M}(0) \) is an invariantly defined region just outside the chronology violating region which is covered by extremely short spacelike geodesics. In a tubelike region along any one of these geodesics the metric can be put in the form

\[ ds^2 = dl^2 + g_{ab}^{(2+1)}(l, x_\perp) \, dx^a_\perp \, dx^b_\perp, \]

subject to the boundary condition

\[ g_{ab}^{(2+1)}(0, 0_\perp) = g_{ab}^{(2+1)}(\ell, 0_\perp). \]

If we now Fourier decompose the metric in this tubelike region the boundary conditions imply that \( p_\ell = n \hbar / \ell \). For \( \ell < L_{\text{Planck}} \), high-momentum trans-Planckian modes \( p_\ell > n \hbar / L_{\text{Planck}} = n E_{\text{Planck}} / c \) are an unavoidable part of the analysis. That is, close enough to the chronology violating region one is intrinsically confronted with Planck scale physics; and
the onset of Planck-scale physics can be invariantly characterized by the length of short but nontrivial spacelike geodesics. In particular the relevant Planck scale physics includes Planck scale fluctuations in the metric — these fluctuations in the geometry of spacetime fuzz out the manifold picture that is the essential backdrop of semiclassical gravity. Thus quantum physics wins the day, and curved space quantum field theory is simply not enough to complete the job.

Overall, this entire chain of development has led the community to a conclusion diametrically opposed to the initial hopes of the early 1990’s — the hopes for a simple and universal classical or semiclassical mechanism leading to chronology protection seem to be dashed, and the relativity community is now faced with the daunting prospect of understanding full quantum gravity just to place notions of global causality on a firm footing.

Where we stand

There is ample evidence that quantum field theory is a good description of reality, and there is also ample evidence that general relativity (Einstein gravity) is a good description of reality. From the obvious statement that in our terrestrial environment gravity is well described by classical general relativity, while condensed matter physics is well described by quantum physics, it follows that semiclassical quantum gravity (curved space quantum field theory with the Einstein equations coupled to the quantum expectation value of the stress-energy) is a more than adequate model over a wide range of situations. (No-one seriously doubts the applicability of semiclassical gravity to planets, stars, galaxies, or even to cosmology itself once the universe has emerged from the Planck era.)

Nevertheless, there are apparently plausible situations in semiclassical gravity that naively seem to lead to the onset of causality violation; and attempts at protecting chronology inevitably lead one back to considerations of full quantum gravity. The situation is somewhat reminiscent of black hole physics where the infinite redshift at the black hole horizon is often interpreted as a microscope that could potentially open a window on the Planck regime [36, 37]. Similarly, in discussing chronology protection the region near the chronology horizon is subject to Planck scale physics (believed to include Planck scale fluctuations in the geometry) so that semiclassical gravity is not a “reliable” guide near the chronology horizon [38, 39]. This opens a second window on Planck scale physics — though the chances of experimentally building a time machine (or getting close enough to forming a chronology horizon to actually see what happens) must be viewed as even somewhat less likely than the chances of experimentally building a general relativity black hole. (Black hole analogues, such as acoustic dumb holes, are another story [40, 41, 42].)
One possible response, given that we will inevitably have to face full-fledged quantum gravity, is to take chronology protection as being so basic a property that we should use it as a guide in developing our theory of quantum gravity:

1. As already mentioned, canonical gravity, whatever its limitations in other areas, does automatically enforce chronology protection by its very construction. Canonical quantum gravity certainly has serious limitations, but it does at least provide a firm kinematic foundation.

2. Lorentzian lattice quantum gravity, as championed by Ambjorn and Loll, also enforces chronology protection by construction [43, 44, 45, 46]. It does so by only summing over a subset of Euclidean lattice geometries, a subset that is compatible with a global Wick rotation back to globally hyperbolic Lorentzian spacetime. At least in low dimensionality, large low-curvature regions of spacetime emerge (large compared to the Planck length, small sub-Planckian curvature). These regions are suitable arenas for curved-space quantum field theory. There are however many loose ends to work out — such as the details of the emergence of the Einstein–Hilbert action in the low-energy limit.

3. Quantum geometry (Ashtekar new variables) is still in a state where details concerning the emergence of a “continuum limit” are far from settled; in particular it is not yet in a position to say anything about chronology protection one way or the other.

4. Brane models (nee string theory) are also not yet able to address this issue. In the low-energy limit brane models are essentially a special case of semiclassical quantum gravity, with the brane physics enforcing a particular choice of low-energy quantum fields on spacetime. In this limit, brane models have nothing additional to say beyond generic semiclassical gravity. In the high-energy limit where the physics becomes “strongly stringy” the entire manifold picture seems to lose its relevance, and there is as yet no reliable formulation of the notion of causality in the string regime. One possibility is to use string dualities: If the strongly-coupled string regime is dual to a weakly-coupled regime where the manifold picture does make sense, then you can at least begin to formulate local notions of causality in the weakly coupled regime and then bootstrap them back to the strongly-coupled regime via duality. But then you still have to decide which class of geometries you will permit in the weakly-coupled regime (globally hyperbolic? stably causal?), and the overall situation is far from clear.
So, is chronology protected? Despite a decade’s work we do not know for certain, but I think it fair to say that the bulk of physicists looking at the issue believe that something along the lines envisaged by Stephen in his “chronology protection conjecture” will ultimately save the day, as Stephen puts it:

*There is also strong experimental evidence in favour of the conjecture — from the fact that we have not been invaded by hordes of tourists from the future.*

It seems to me that approaches based on Novikov’s consistency condition [24, 25, 26] are now somewhat in disfavour, largely on philosophical rather than physical grounds. The same comment applies to attempts at invoking the many-worlds interpretation of quantum physics, or other ways of radially re-writing the foundations of physics. Still, despite their relative unpopularity (or maybe, because of their relative unpopularity) these more radical alternatives should also be kept in mind as exploration continues. Unfortunately, if chronology protection is the answer, we will have to wander deep into the guts of quantum gravity to know for certain.
References


References


