

Chronological structure of a Gödel type universe with negative cosmological constant

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ABSTRACT

We show that the Gödel solution of five-dimensional gauged supergravity contains either closed time-like curves through every space-time point or none at all, dependent on the rotational parameter. In addition, we present a deformation of that solution with a parameter κ which characterizes the symmetry of four-dimensional base space: for $\kappa = 1, 0, -1$ it has spherical, flat and hyperbolic symmetry, respectively. Also investigated are the causal properties of the lifted solution in 10 dimensions.

The occurrence of closed time-like curves (CTCs) is widely viewed as a pathological feature of space-time. While there certainly are solutions of general relativity that contain CTCs (a review of the classical examples can be found in [1]), there is an on-going discussion whether or not a more complete description, in particular one that includes quantum effects, would lead to the absence of CTCs [2, 3, 4]. One widely-held view is that a viable theory of quantum gravity should not permit their existence, at least not as far as curves accessible to an outside observer are concerned (as opposed to, say, the CTCs hidden behind the horizon of a rotating black hole). However, in recent time it was realized that string theory, as one candidate for a theory of quantum gravity, does allow for vacua with “naked” CTCs, and in particular that some such vacua preserve at least one supersymmetry and thus can be expected to be reasonably stable.

The first such example was the rotating black hole solution in 5 dimensions [5, 6] where, for sufficiently large angular momentum, CTCs appear *outside* the horizon [7, 8, 9]. If one embeds such black hole into an asymptotic Anti-de Sitter space [10], it acquires an ergosphere which contains CTCs for arbitrary values of the angular momentum [11]. Another class of string solutions are supertubes, where the space-time contains CTCs in the over-rotating case [12], analogous to the asymptotically flat black holes. There is an on-going discussion about whether one should exclude these vacua in string theory, or at least show that the pathological regions do no harm – either since they somehow cannot be reached by geodesic motion or because a more fully string-theoretical description indicates the need to excise and to replace them with non-pathological space-time patches. More specifically, a typical test involves the examination of these backgrounds using different kinds of brane probes and investigates these probes’ world volume theory. Indications that the probes cannot reach the pathological regions (behind the “CTC horizon”) include the appearance of ghost degrees of freedom (with negative kinetic energies) in the world-volume theory, see for example [8, 13]. Another approach has examined the question of whether a problematic black-hole space-time containing CTCs can be created in some way from CTC-free space-times, using the string-theoretical microscopic description for the black hole in question [14]. However, while certainly valuable, arguments of this type are limited insofar as they do not address “eternally existing” pathological configurations. (Another possible problem with the argument in the reference cited is its reliance on an energy condition that might be violated in quantum gravity).

Yet another class of chronology violating solutions, and one which has attracted much attention recently, is the Gödel-type solution of 5-dimensional ungauged supergravity, where the matter of the original Gödel solution is replaced by a magnetic configuration of the Abelian gauge field [15]. This homogeneous solution describes a rotating universe and has CTCs through every point of space-time. A number of interesting properties of this solution has been explored: The CTCs can be shielded by holographic screens, and the solution is T-dual to certain pp-waves [16, 17]; black holes can be embedded into this space-time [8, 18, 19, 20, 21], and it has been examined using various probes

[22, 23]. Another interesting aspect dates back to the original discovery of the solution, which can be obtained by dimensional reduction of a WZW model discussed by Nappi and Witten [24] that is itself a special case of a larger class of conformal field theories introduced in [25]. Interestingly enough, these CFTs can be solved exactly [26]. The fact that the partition function of this background becomes divergent has been linked to the appearance of CTCs and interpreted as strings becoming tensionless [27], however, a physical understanding of this phase transition is still missing. Let us also mention that there is an interpolating string solution between this solution and the exact flux backgrounds of the Melvin type [28], which exhibits similar instabilities [29].

In [30] the corresponding Gödel solution of minimal gauged supergravity was presented.¹ This solution also exhibits CTCs, but these have not so far been discussed in detail. It solves the Einstein-Maxwell equations associated with the bosonic Lagrangian of minimal five-dimensional supergravity,

$$*\mathcal{L} \sim *1 \left[\frac{R}{2} + \chi^2 - \frac{1}{4}F^2 \right] - \frac{2}{3\sqrt{6}} F \wedge F \wedge A , \quad (1)$$

in which the cosmological constant χ^2 appears due to a Abelian gauging of the R-symmetry $SU(2)$. By construction within the framework of [30], the solution is supersymmetric. It belongs to a class of solutions defined by a stationary metric with a Kähler base space; preservation of supersymmetry entails that the field strength of the Abelian gauge field be anti-selfdual over the Kähler space, and have no components in the time-direction. In the form discussed in [30], the Gödel-type solution is based on the Bergmann metric parameterizing the coset manifold $\frac{SU(2,1)}{U(2)}$, its five-dimensional metric and gauge field are²

$$ds^2 = -(dt + \omega)^2 + \frac{d\varrho^2}{4\varrho V(\varrho)} + \frac{\varrho}{4} \left[V(\varrho)\sigma_3^2 + \sigma_1^2 + \sigma_2^2 \right] \quad \text{and} \quad A = \frac{\sqrt{3}f\varrho}{8V(\varrho)} \sigma_1 \quad (2)$$

where

$$\omega = \frac{\chi\varrho}{2\sqrt{6}} \sigma_3 + \frac{f\varrho}{4\sqrt{2}V(\varrho)} \sigma_1 \quad , \quad V(\varrho) := 1 + \frac{\chi^2}{6}\varrho \quad (3)$$

and where the left-invariant $SU(2)$ 1-forms

$$\begin{aligned} \sigma_1 &= \sin \phi d\theta - \cos \phi \sin \theta d\psi , \\ \sigma_2 &= \cos \phi d\theta + \sin \phi \sin \theta d\psi , \\ \sigma_3 &= d\phi + \cos \theta d\psi \end{aligned} \quad (4)$$

¹The four-dimensional cousin of this space-time was recently found in [31].

²Compared with the form given in [30], we have rescaled field strength and cosmological constant, chosen the (irrelevant) phase angle between \mathcal{F}_1 and \mathcal{F}_2 to be zero, and performed a coordinate redefinition $\varrho = 6/\chi^2 \sinh^2(\chi r/\sqrt{6})$ to remove cumbersome hyperbolic functions.

were used in order to specify our solution in terms of Eulerian coordinates $0 \leq \rho, 0 \leq \theta \leq \pi, 0 \leq \psi \leq 2\pi$ and $0 \leq \phi < 4\pi$. In the limit of vanishing cosmological constant, $\chi \rightarrow 0$, one recovers the supersymmetric Gödel solution, which, in Cartesian coordinates, reads [15]

$$\begin{aligned} ds^2 &= -(dt + \omega)^2 + dx^m dx^m \\ \omega &= \frac{2}{\sqrt{3}}A = 2\gamma(x^1 dx^2 - x^2 dx^1 + x^3 dx^4 - x^4 dx^3) . \end{aligned} \quad (5)$$

Regarding the $\chi = 0$ solution, the proof that there are CTCs through every point is greatly helped by the fact that the space-time is homogeneous. Once the existence of any CTC, however special, has been proven, homogeneity ensures that CTCs exist everywhere. This argument does not carry over to the Gödel solution in gauged supergravity, however, where homogeneity is lost. So far, it has not yet been shown whether there are regions of this space-time that are free of CTCs. As we will discuss in detail now, this is unfortunately not the case.

Before this discussion, however, let us present a brief collection of further information about the gauged five-dimensional Gödel universe. First of all, there is an easy generalization of the solution in the following way: For the base space $\frac{SU(2,1)}{U(2)}$, different parametrizations are possible. In the case of the Bergmann metric, the base metric is a complex line bundle over S_2 ; as a generalization, one can admit a complex line bundle over other two-dimensional spaces of constant curvature κ . This corresponds to the $n \rightarrow \infty$ limit of the metric described in [32] and from this base-space five-dimensional solutions are readily constructed, namely

$$\begin{aligned} ds^2 &= -(dt + \omega)^2 + \frac{d\varrho^2}{4\varrho V(\varrho)} + \frac{\varrho}{4}V(\varrho) \left[d\phi' + \frac{1}{2} \frac{xdy - ydx}{1 + \frac{\kappa}{4}(x^2 + y^2)} \right]^2 + \frac{\varrho}{4} \frac{dx^2 + dy^2}{[1 + \frac{\kappa}{4}(x^2 + y^2)]^2} \\ A &= \frac{\sqrt{3}f}{8} \left[\frac{\varrho}{\kappa V(\varrho)[1 + \frac{\kappa}{4}(x^2 + y^2)]} (\sin \kappa\phi' dx + \cos \kappa\phi' dy) - \frac{6}{\kappa\chi^2} dy \right], \end{aligned} \quad (6)$$

where κ can be scaled so that three different cases $\kappa = -1, 0, 1$ remain. In this definition,

$$\omega = \frac{\chi\varrho}{2\sqrt{6}} \left[d\phi' + \frac{1}{2} \frac{xdy - ydx}{1 + \frac{\kappa}{4}(x^2 + y^2)} \right] + A \quad , \quad V(\varrho) = \kappa + \frac{\chi^2}{6} \varrho . \quad (7)$$

(Note that we have added a constant part to A_y to ensure a smooth limit $\kappa = 0$.) In order to avoid a conical singularity, one should require a periodicity $\phi' \simeq \phi' + \frac{4\pi}{\kappa}$; in particular, for $\kappa = 0$ all coordinates are non-compact. Let us note here another possible choice for the gauge connection, namely

$$A_x + iA_y = -\frac{1}{\kappa} \log \left[\left(\frac{6\kappa}{\chi^2\varrho} + 1 \right) \left(1 + \frac{\kappa}{4}[x^2 + y^2] \right) e^{i\phi'} \right] (dx + idy) .$$

The ambiguity in the choice of gauge connection can be understood as follows: The (complex) field strength two-form is conformally equivalent to the two non-exact two-forms of the quaternionic space $\frac{SU(2,1)}{U(2)}$, with the conformal factor defined by the Kähler

potential. However, the split into an exact two-form and a pre-factor is not unique, and it might be interesting to explore the possible relation of this to Kähler transformations.

Having noted the existence of this one-parameter family of solutions, we now focus on the spherical case $\kappa = 1$. It corresponds to the Gödel-type solution discussed before, and the metric (2) can be recovered changing coordinates as $x + iy = 2 e^{i\psi} \tan \theta/2$.

For the discussion of CTCs, the best starting point are left- and right-invariant vectors associated with the Euler angle coordinates – natural candidates for tangent fields of congruences of closed curves. In flat space with standard metric, these are all Killing vectors; constructing the space-time with arbitrary values of χ and f , however, most of them lose their Killing property (of course, other Killing vectors are gained, notably $\xi_0 = \partial_t$ associated with the stationarity of the space-time). First of all, there are the $SU(2)$ -right-invariant vectors

$$\begin{aligned}\xi_1^R &= -\sin \psi \partial_\theta - \cot \theta \cos \psi \partial_\psi + \frac{\cos \psi}{\sin \theta} \partial_\phi \\ \xi_2^R &= \cos \psi \partial_\theta - \cot \theta \sin \psi \partial_\psi + \frac{\sin \psi}{\sin \theta} \partial_\phi \\ \xi_3^R &= \partial_\psi.\end{aligned}\tag{8}$$

These all remain Killing vectors of the full space-time. Secondly, there are the $SU(2)$ -left-invariant vectors

$$\begin{aligned}\xi_1^L &= -\frac{\cos \phi}{\sin \theta} \partial_\psi + \sin \phi \partial_\theta + \cot \theta \cos \phi \partial_\phi \\ \xi_2^L &= \cos \phi \partial_\theta + \frac{\sin \phi}{\sin \theta} \partial_\psi - \cot \theta \sin \phi \partial_\phi \\ \xi_3^L &= \partial_\phi.\end{aligned}\tag{9}$$

Dual as they are to the left-invariant forms used in the construction of the metric, the products between these vectors are comparatively simple, and we will thus use these vectors below in the construction of CTCs. They are, however, Killing only in special limits: ξ_3^L becomes a Killing vector in the limit $f = 0$, for arbitrary χ , when the gauge field is switched off, the Einstein equations reduce to $Ric = -\frac{2}{3}\chi^2 g$ and the space-time becomes Anti-deSitter in unusual coordinates (to return to static coordinates, one need merely shift $\phi \rightarrow \phi' = \phi - \frac{2\chi}{\sqrt{6}}t$). On the other hand, ξ_1^L becomes a Killing vector in the limit $\chi = 0$ of the non-gauged Gödel solution.

With these preparations, now for an analysis of the possible CTCs in the space-time (2). As it turns out, the presence or absence of CTCs in that space-time is governed by the relative values of the parameters χ and f : As long as

$$f^2 \leq \frac{4}{3} \chi^2\tag{10}$$

there are no CTCs; otherwise, there are CTCs through every space-time point. (Note

that this is a more stringent limit than the one that resulted from the construction of special CTCs in [30]; the limit given there would, in our conventions³ be $f^2 \leq 16/3\chi^2$.)

This can be seen as follows. First, we can ask ourselves under what conditions the coordinates chosen in (2) define a global foliation of the space-time they describe, in other words: under what conditions is the hypersurface Σ_t , defined by ρ, θ, ψ, ϕ at $t = \text{const.}$, spacelike everywhere? The best way to find out is a look at the eigenvalues of the induced metric on each such hypersurface. Doing so (a basis of $\partial_t, \partial_\rho, \xi_1^L, \xi_2^L, \xi_3^L$ proves convenient), we see that it has three eigenvalues that are positive at any ρ for arbitrary values of f and χ (zeroes at $\rho = 0$ turn out to be coordinate artefacts, as can be seen by changing to cartesian coordinates). Not so for the fourth eigenvalue: As long as $f^2 > 4/3\chi^2$, it starts out as positive for small ρ , has a zero at $\rho_0 = 24/(3f^2 - 4\chi^2)$ and is negative for all larger values of ρ . This tells us directly that, as long as $f^2 \leq \frac{4}{3}\chi^2$ all is well, and the space-time contains no CTCs. Viewing this as a bound on the parameter f^2 , which measures the angular momentum/the magnetic flux, the situation is somewhat similar to the rotating black holes, with the presence of CTCs governed by angular momentum versus total mass.

Next for the proof that, as long as $f^2 > 4/3\chi^2$, there are CTCs through *every* space-time point. First, let us look at CTCs that remain at constant coordinate distance $\rho = \text{const.}$ from the origin. In order to find such curves, it is no longer sufficient to consider only integral curves along one suitably chosen angular coordinate, as was possible in the non-gauged case [15] or in the case of supersymmetric, five-dimensional black holes [11]. Instead, a more general Ansatz is needed. Our above analysis of the induced metric on the hypersurfaces Σ_t suggests a likely candidate for a tangent vector field build from the ξ_i^L , namely the eigenvector associated with the eigenvalue that is responsible for the degeneracy of the metric of Σ_t at $\rho = \rho_0$,

$$\xi_c = \left[\sqrt{1 + \frac{3f^2}{16\chi^2 V(\rho)^2} - \frac{\sqrt{3}f}{4\chi V(\rho)}} \right] \xi_3^L + \xi_1^L. \quad (11)$$

It is readily checked that $|\xi_c|^2 < 0$ for $\rho \geq \rho_0$, as long as $f^2 > 4/3\chi^2$. It still needs to be shown that the congruence defined by the tangent field ξ_c consists of *closed* curves. This can be seen as follows: As manifolds, the surfaces $\rho = \text{const.}$ are homeomorphic to three-spheres. If any such three-sphere is equipped with the standard metric, then the ξ_i^L are Killing vectors, and the geodesics are great circles. In particular, the projection of the tangent vector onto any ξ_i^L is constant along each geodesic so, as the ξ_i^L form a basis field for the tangent space of the S^3 , each geodesic is an integral curve of some linear combination of the ξ_i^L and, conversely, the integral curve of any such linear combination is closed.

So far, we have shown that, for suitable f , CTCs exist in some regions of space, namely for $\rho > \rho_0$. The fact that CTCs pass through all points of the region $\rho \leq \rho_0$, as well is plausible, as those closed timelike curves whose existence we have shown should

³Notably $\mathcal{F}_1^2 + \mathcal{F}_2^2 = f^2/2$ and $\chi_{\text{there}} = \sqrt{2} \cdot \chi_{\text{here}}$.

be easily deformable to “timelike winding stairs”, i.e. open curves that lead from a space-time point $(t, \rho, \theta, \psi, \phi)$ to $(t - \Delta t, \rho, \theta, \psi, \phi)$ for some finite Δt . Given such spiral curves, we can travel from $\rho < \rho_0$ into the region $\rho > \rho_0$ in an ordinary time-like curve, descend the spiral to some suitably earlier coordinate time, and return to our initial ρ value and our initial coordinate time. More concretely, taking as one of a plethora of possibilities the Ansatz⁴

$$\xi_{ws} = \frac{2\chi}{\sqrt{3}f} \left[1 + \frac{(3f^2 - 4\chi^2)^2}{9f^2(3f^2 + 4\chi^2)} \right] \xi_3^L + \xi_1^L - c \partial_t. \quad (12)$$

With this definition, $|\xi_{ws}|^2$ is a quadratic function of the as yet unfixed coefficient c with two zeroes c_{\pm} . Both for $c \ll 0$ and for $c \gg 0$ we find $|\xi_{ws}|^2 < 0$ as expected: choosing $c \ll 0$, we will end up near-parallel to $-\partial_t$, and thus in the past light-cone; choosing $c \gg 0$, near-parallel to ∂_t and in the future light-cone. The necessary and sufficient condition for ξ_{ws} to be tangent to a winding stair curve is for the smaller zero, c_- , to be *positive*. Calculation of $|\xi_{ws}|^2$ at, for instance, $\rho = 2\rho_0$ shows this to be the case as long as $f^2 > 4/3\chi^2$, and choosing $c = c_-/2 > 0$ defines an example of a vector field tangent to a congruence of winding-stair curves. This completes the proof of our statement that every point in space-time is on a CTC as long the cosmological constant is small, but in contrast to the case of vanishing cosmological constant all CTC disappear in the moment the bound (10) is saturated.

Finally, it is of interest to lift the five-dimensional to string-friendly ten dimensions; in the case of black holes or the Gödel solution without cosmological constant, this lifting offers a way to get rid of the offending CTCs [33, 8]. More concretely, we lift to a solution of ten-dimensional type IIB-supergravity following the prescription given in section II of [34]. Let the additional five (angular) coordinates be α, β and $\phi_i, i = 1, 2, 3$, and define functions μ_i of α, β such that⁵ $\sum_{i=1}^3 \mu_i^2 = 1$. The lifted metric constructed from five-dimensional line element ds_5^2 and the one-form A is

$$ds_{10}^2 = ds_5^2 + \frac{6}{\chi^2} \sum_{i=1}^3 \left[(d\mu_i)^2 + \mu_i^2 \left(d\phi_i + \frac{\chi}{3} A \right)^2 \right] \quad (13)$$

and, together with the five-form $\bar{F} = G + *G$ defined by

$$G = \frac{1}{2\sqrt{5}} \left[\frac{\chi}{3} \varepsilon^{(5)} - \frac{1}{2\chi^2} \sum_{i=1}^3 d(\mu_i^2) \wedge d\phi_i \wedge *^{(5)} F \right], \quad (14)$$

where $\varepsilon^{(5)}$ and $*^{(5)}$ are the five-dimensional volume form and Hodge dualization, respectively. It fulfills the ten-dimensional Einstein-Maxwell-equations

$$Ric_{MN}^{(10)} - 5\bar{F}_{MM_2\dots M_5} \bar{F}_N{}^{M_2\dots M_5} = 0 \quad \text{and} \quad d * F = 0. \quad (15)$$

⁴This is the above eigenvector ϵ_c , with the coefficient of ξ_3^L expanded around ρ_0 and evaluated at an arbitrarily chosen point $\rho = \rho_0 + 2/\chi^2$. It has the advantage of being simple enough for $|\xi_{ws}|^2 < 0$ to be demonstrable with modest effort.

⁵For concreteness, one can choose $\mu_1 = \sin \alpha, \mu_2 = \cos \alpha \sin \beta, \mu_3 = \cos \alpha \cos \beta$.

One might have hoped that this lifting alone could cure the illness of the CTCs, as it involves a shift in the $U(1)$ fibre that is at the root of the five-dimensional solution's chronological problems. However, this is not the case. Again, by looking at the foliation defined by the given coordinates, it follows that for $f^2 \leq 4/3\chi^2$, there are no CTCs, and again, for $f^2 > 4/3\chi^2$, CTCs through every space-time point can be constructed as follows, in perfect analogy to the procedure we have used in five dimensions: To show the existence of CTCs through all points in the region $\rho > \rho_0^{(10)} = 12/(\chi[\sqrt{3}f - 2\chi])$, one can use the tangent vector field $\xi_3^L + \xi_1^L$. For the remaining regions, study of the tangent field $-c\partial_t + \xi_3^L + \xi_1^L$ shows for $f^2 \leq 4/3\chi^2$ and some $\rho_{ws} > \rho_0$ the existence of “winding stair” curves that can be used to go back in time and hence to construct CTCs of finite length through space-time points with $\rho < \rho_0$.

There are a number of open questions which we were not able to discuss in detail in this note, but which might be useful for future work. Notably, the Gödel solution we examined can be seen as a deformation of AdS_5 and, as we have shown, CTCs exist even close to the boundary. This raises the question of what happens in the dual field theory. As we have mentioned, for the Gödel solution in ungauged supergravity there exist a description in string theory in terms of a solvable CFT; it is natural to ask whether the solutions of gauged supergravity discussed here can be described by an exact model.

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