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Black hole formation from collisions of cosmic strings

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Abstract

We explain simple semi-classical rules to estimate the lifetime of any given highly excited quantum state of the string spectrum and apply them to identify new long-lived string states. Using analytic formulae for the string evolution after joining and interconnection, we study examples of fundamental cosmic string collisions leading to gravitational collapse. We find that the interconnection of two strings of equal and opposite maximal angular momenta and arbitrarily large mass generically leads to the formation of black holes. (Based on the works (Iengo and Russo 2006 *J. High Energy Phys.* JHEP02(2006)041, Iengo and Russo 2006 *J. High Energy Phys.* JHEP08(2006)079).)

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1. Introduction

Recent works have indicated that in brane inflation models cosmic strings are copiously produced during the brane collision [3, 4]. This has led to a renewed interest in the physics of cosmic strings and to consider the exciting possibility that there could be long-lived fundamental strings of cosmic size (for reviews and more references, see [5, 6]). Finding such objects could constitute a test of string theory.

The interpretation of a cosmological strings in terms of massive fundamental strings faces a basic problem, pointed out long ago in [7]: generic massive string states are expected to be very unstable. It is important to understand which string models can give rise to long-lived string states. Computing the lifetime of a given massive string state is extremely complicated, due to the large degeneracy of the decay products. It has been done only in few cases for especially simple states (see [1, 8–11]). The different calculations performed in these works give us an insight on what are the dominant decay channels for massive string states and how the decay rate depends on the mass and on the 'shape' of the string. In particular, one finds that, surprisingly, there are strings which become more stable for larger masses [10, 11].

Combining the different known results, in the first part of this work we will first derive simple rules to make a good estimate of the decay rate by splitting and massless radiation, and hence the lifetime of an arbitrary string state. A direct application of these results will be the construction of new long-lived string states that can be used in cosmological models.

The dynamics of cosmic strings could also lead to interesting astrophysical events such as gravitational waves or black hole formation that could result from the evolution of a single cosmic string [12–14] or from the collision of cosmic strings. In the second part of this work we will study the latter problem by using the formalism of string splitting, joining and intercommutation developed in [1]. Surprisingly, we will find that gravitational collapse is a quite common phenomenon ensuing the encounter of strings of equal and opposite maximal angular momenta.

2. Decay rate due to breaking

In general, a string can decay either by emitting light particles or by splitting into two massive strings. Which channel is dominant depends on the string and on the number of uncompact dimensions. For splitting into two massive strings with large masses, the two outgoing strings are highly excited string states which admit a classical description. The decay rate in this case can be estimated by semiclassical arguments.

2.1. Open strings

The probability of breaking for an open string was studied in [15]. There it was suggested that it is constant along the string and proportional to g_o^2. We propose that, more precisely, the probability of breaking on a given point once in a period is

$$P_o = \frac{g_o^2}{T}, \tag{1}$$

where T is the period of oscillation (so the probability of breaking at any moment is $T P_o = g_o^2$ in agreement with [15]). In the gauge $X_0 = \frac{T}{\pi}\tau$, we have $T = 2\pi\alpha' M$, where M is the string mass. The decay rate for an open string is thus obtained by multiplying P_o by the number of points of the string L/l_s, where $L \sim \alpha' M$ is the length and $l_s = \sqrt{\alpha'}$, and by the number of 'instants' in one period T/l_s. We find

$$\Gamma_{\text{open}} \cong \left(\frac{L}{l_s}\right)\left(\frac{T}{l_s}\right) P_o \cong g_o^2 \frac{L}{l_s^2}. \tag{2}$$

This formula precisely reproduces the law $\Gamma_{\text{open}} \cong g_o^2 M$ derived in [8] by the full quantum calculation of the decay rate for the open string with maximum angular momentum.

2.2. Closed strings

For a closed string, in the absence of D-branes, breaking is possible only if two points of the string meet. There are several possible configurations, as discussed below. In particular, if the distance between any two points of the string is $\gg l_s$ at all times, then the string is 'unbreakable', meaning that the decay rate into two massive $\left(M_{1,2} \gg l_s^{-1}\right)$ string states is exponentially suppressed, $\Gamma = O\left(\mathrm{e}^{-cM^2}\right)$.

2.2.1. Folded string. The two folds of the closed string are in permanent contact, so the string can break at any time. We can estimate the decay rate by viewing it as two open strings on the top of each other. The breaking can take place only if at a given time the two open strings break at the same point, up to an uncertainty of order l_s. The decay rate is thus

$$\Gamma_{\text{folded}} \cong \left(\frac{L}{l_s}\right)\left(\frac{T}{l_s}\right) P_o P_o l_s \cong g_s^2 \frac{L}{l_s^3 M} \cong g_s^2 l_s^{-1}, \qquad g_s = g_o^2. \tag{3}$$

Remarkably, unlike the open string, it is constant independent of M. This semiclassical estimate precisely reproduces the quantum decay rate obtained in [1, 10] from the exact evaluation of $\text{Im}(\Delta M^2)$ at one loop.

2.2.2. String which becomes folded at an instant of time. An explicit example is the squashing ellipse of [1]. The classical string solution is given by

$$\begin{aligned} x_1 &= 2L\cos\theta\cos\tau\cos\sigma, & x_2 &= 2L\sin\tau\cos\sigma, \\ x_3 &= 2L\sin\theta\sin\tau\sin\sigma, & x_0 &= 2L\tau, \end{aligned} \tag{4}$$

where θ is a parameter and $\sigma \in [0, 2\pi)$. For θ generic, it describes an ellipse that rotates around one of its axes and simultaneously performs pulsations, with the one of its radii (the one on the axes of rotation) becoming zero at $\tau = n\pi$, $n =$ integer. This string interpolates between the folded string ($\theta = 0$) and the pulsating circular string ($\theta = \pi/2$). At each period of oscillation, there are two times where the string (4) becomes folded and it can break. Quantum mechanically, the breaking process is important during the time that the smaller radius of the ellipse has size $<l_s$. This occurs during a time $\Delta x_0 \cong l_s$ at each period. So the fraction of time where the string can break is l_s/L. By definition, the decay rate is the number of events at each period of oscillation of the string. This means that the decay rate will be given by

$$\Gamma_{\text{squash}} \cong \frac{l_s}{L}\Gamma_{\text{folded}} \cong g_s^2 \frac{1}{l_s\sqrt{N}}. \tag{5}$$

This is in precise agreement with the explicit calculation of the quantum decay performed in [1], providing a non-trivial confirmation of the semiclassical rules given above.

2.2.3. Pulsating circular string. This is a circular string which shrinks to a point once in a period. The decay rate can be calculated just as in the squashing string case, but now taking into account that the two points where the breaking takes place are arbitrary, since by the time the string is completely shrunk all points are in contact. This means that there is an additional factor of L/l_s. Thus we get

$$\Gamma_{\text{pulsating}} \cong g_s^2 \frac{L^2}{M^2 l_s^5} \cong g_s^2 l_s^{-1}, \tag{6}$$

where we have used that $L \cong l_s^2 M$ for the pulsating circular string. Thus the decay rate is the same as in the folded string case with maximum angular momentum. We have verified this remarkable fact by the exact quantum calculation [1].

3. Decay rate due to massless emission

The quantum massless emission from a closed string contains contributions from four sectors: NS–NS, R–NS, NS–R and R–R. The explicit calculation for every channel has been carried out in the following cases: (a) the string with maximum angular momentum Jmax [10], (b) the rotating ring [11] and (c) the squashing string [1].

The rotating ring is a rigid circular string rotating in two orthogonal planes. In this case, we found [11] that NS–NS emission (which includes graviton, dilaton and antisymmetric tensor) is dominant, whereas R–NS, NS–R and R–R emissions are suppressed by factors $1/N$. Moreover, in this case the NS–NS emission can be accurately described as a radiation process from a classical antenna represented by the classical rotating ring solution. The classical radiation from a source $T_{\mu\nu}(x_0, \vec{x})$ in D uncompact spacetime dimensions is given by

$$\Gamma = g_s^2 \frac{\omega^{D-3}}{M^2} \int \mathrm{d}^{D-2}\Omega \sum_{\xi,\tilde{\xi}} |J|^2, \qquad J = \int \mathrm{d}x_0\, \mathrm{d}\vec{x}\, \mathrm{e}^{\mathrm{i}\omega X_0 - \mathrm{i}\vec{p}.\vec{X}} \xi^\mu \tilde{\xi}^\nu T_{\mu\nu}(X_0, \vec{X}). \tag{7}$$

For a classical string solution, the energy–momentum tensor is

$$T_{\mu\nu} = \int \mathrm{d}\sigma\, \mathrm{d}\tau\, \delta^{(\mathrm{d})}(x - X(\tau, \sigma)) \partial X_\mu \bar{\partial} X_\nu, \tag{8}$$

so that $|J|^2 = |J_R|^2 |J_L|^2$ with (gauge $\xi^0 = \vec{\xi} \cdot \vec{p} = 0$)

$$J_R = \int_0^{2\pi} \mathrm{d}\sigma\, \mathrm{e}^{\mathrm{i}p_- X_{R+}} \vec{\xi} \cdot \partial \vec{X}_R^T(\sigma), \qquad J_L = \int_0^{2\pi} \mathrm{d}\sigma\, \mathrm{e}^{\mathrm{i}p_- X_{L+}} \vec{\xi} \cdot \partial \vec{X}_L^T(\sigma), \tag{9}$$

where we have chosen the frame where the momentum of the emitted massless particle is $p^\mu = (\omega, -\omega, \vec{0})$ and the gauge $X_0 = \alpha' M\tau = 2\sqrt{\alpha' N}\tau$. $X_\pm$ refer to the light-cone coordinates, where $p_+ = 0$. The radiated energy is $\omega = \frac{M^2 - M'^2}{2M} = N_0/(2\sqrt{N})$, where we have set $\alpha' = 4$ and $N_0 \equiv N - N'$, with $M = \sqrt{N}$ being the mass of the original state and $M' = \sqrt{N'}$ being the mass of the massive state after the emission.

The classical formula is expected to hold for $\omega \ll O(1/\sqrt{\alpha'})$, i.e. for $N_0 \ll \sqrt{N}$. If the massless NS–NS emission with higher energies is suppressed, then the classical formula can be used to compute the total radiation emission. In general, since $\mathrm{i}p_- X_{R,L+}(\sigma) \sim \mathrm{i}N_0 f_{R,L}(\sigma)$ in the exponent in (9), then $J_{R,L}$ are exponentially suppressed as a function of N_0, unless there is a saddle point in the integration over σ in (9), or a 'kink' in the function $f_{R,L}(\sigma)$. A saddle point ('cusp') occurs if $\partial_\sigma X_{R+}$ and $\partial_\sigma X_{L+}$ vanish for some σ, while a kink occurs when there is discontinuity in the first derivative in the function $f_{R,L}(\sigma)$. This is the case of the cusp and kink string configurations studied in [12, 16, 17]. So let us first assume that there is no cusp or kink. In this case, $J_{R,L} = M h_{R,L}(N_0, \Omega)$, where $h_{R,L}(N_0, \Omega)$ are exponentially suppressed for $N_0 \gg 1$. We have used the fact that $\vec{\xi} \cdot \partial \vec{X}_{R,L}^T$ is generically proportional to M as can be seen from the Virasoro constraints. Therefore, the decay rate (7) is given by

$$\Gamma \cong \mathrm{const}\, g_s^2 M^{5-D} N_0^{D-3} \int \mathrm{d}^{D-2}\Omega |h_R(N_0, \Omega) h_L(N_0, \Omega)|^2. \tag{10}$$

The total rate is obtained by summing over N_0 (i.e. all possible energies of the massless particle). This sum is convergent and one obtains

$$\Gamma_{\text{total rad}} \cong \mathrm{const}\, g_s^2 M^{5-D}. \tag{11}$$

This law has been verified explicitly in [11] by both the classical and quantum calculation for the rotating ring, which has no cusp or kink. In this case, since decays into massive channels are exponentially suppressed, $\Gamma_{\text{total rad}} \cong \Gamma_{\text{total}}$, so the lifetime of the ring state is $\sim g_s^{-2} M^{D-5}$. Note that the time for the state to radiate away all of its energy is much longer by a factor proportional to $N = M^2$.

Let us now consider string configurations with cusps. Whenever the sum over N_0 is convergent for every angle one obtains again, either in the cusp or kink case, $\Gamma_{\text{total rad}} \cong \mathrm{const}\, g_s^2 M^{5-D}$. By comparing with the exact numerical results of the full quantum computation, we have verified that this is indeed the behaviour in $D = 4$ for Jmax and the squashing string.

Finally, another case where the massless radiation emission can be computed explicitly is that of an average string state. In this case, one finds [11, 18]

$$\bar{\Gamma}_{\text{total rad}} \cong g_s^2 M. \tag{12}$$

This formula is the same even if some dimensions are compact.

4. New examples of long-lived cosmic strings

4.1. Rotating straight string on $M^4 \times S^1$

Let t, X, Y, Z represent the uncompact coordinates of M^4 ((3+1)-dimensional Minkowski space) and W compact dimensions of radius R. The solution is as follows:

$$\begin{aligned}
&X = L\cos\tau\cos\sigma, && X_R(\sigma_-) = \tfrac{1}{2}L\cos\sigma_-, && X_L(\sigma_+) = \tfrac{1}{2}L\cos\sigma_+,\\
&Y = L\sin\tau\cos\sigma, && Y_R(\sigma_-) = -\tfrac{1}{2}L\sin\sigma_-, && Y_L(\sigma_+) = \tfrac{1}{2}L\sin\sigma_+,\\
&W = nR\sigma, && W_R(\sigma_-) = \tfrac{1}{2}nR\sigma_-, && W_L(\sigma_+) = \tfrac{1}{2}nR\sigma_+,\\
&t = \kappa\tau, && \kappa = \sqrt{L^2 + n^2R^2}, &&
\end{aligned} \tag{13}$$

where $\sigma_\pm = \sigma \pm \tau, \sigma \in [0, 2\pi)$ and n is an integer representing the winding number. The solution is classically unbreakable for $n = 1$. Although the string looks folded in 3+1 dimensions—and in fact it looks identical to the unstable rotating string of maximal angular momentum [1, 10]—it cannot break because all the points of the string are separated in the internal dimension W. If $R \gg \sqrt{\alpha'}$, then breaking by quantum effects is also suppressed. It can decay by radiation, with a rate (in four dimensions) $\Gamma \sim g_s^2 M, M \sim \mu L$, where $\mu = (2\pi\alpha')^{-1}$ is the string tension and g_s is the closed string coupling constant. The radiation is dominated by soft modes with emitted energy $\omega \sim 1/L$. Thus

$$-\frac{\mathrm{d}M}{\mathrm{d}t} \sim \Gamma \times \omega \cong c_0 g_s^2 \mu, \tag{14}$$

where c_0 is a numerical constant of order 1. Therefore the string takes a time $\sim M/g_s^2$ (or $\sim L/g_s^2$) to substantially decrease its mass.

4.2. Rotating open string which oscillates in extra dimensions

Consider a brane-world model, with a D3-brane placed in the three uncompact directions of our universe. Let x_3 stand for an extra dimension. The open string solution with Dirichlet boundary conditions at $x_3 = 0$ is given by equation (4), but here $\sigma \in [0, \pi)$ whereas for the closed string $\sigma \in [0, 2\pi)$. The solution represents a string rotating in the plane x_1, x_2, with the ends attached to the brane $x_3 = 0$, which at the same time oscillates in the extra dimension x_3. The string can break only at the special times where it lies on the brane, namely $\tau = n\pi$. Using the rules of section 2, one can see that the decay rate by splitting is suppressed for large L. The string nevertheless loses energy at all times by gravitational radiation. For $D = 4$, by summing over N_0 we find that $\Gamma_{\text{graviton}} \sim g_s^2\sqrt{N}$. The lifetime required for a substantial decrease of the energy is again of order $M\big/g_s^2$.

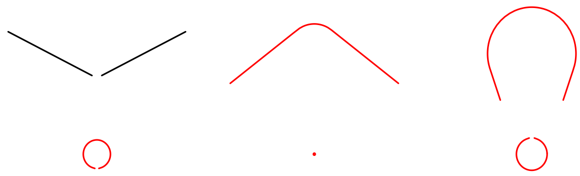

Figure 1. Evolution of the open string which results from the joining of two open strings with maximum and opposite angular momenta.

5. Black hole formation

5.1. By string joining

First, consider two open strings with maximum angular momentum, zero linear momentum and equal energies described by the solutions: $X_{\mathrm{I,II}}(\tau,\sigma) = X_{\mathrm{I,II}L}(\tau+\sigma) + X_{\mathrm{I,II}L}(\tau-\sigma)$ with

$$\begin{aligned} X_{\mathrm{I}}(\sigma,\tau) &= L\cos\sigma\cos\tau, & Y_{\mathrm{I}}(\sigma,\tau) &= L\cos\sigma\sin\tau, \\ X_{\mathrm{I}L}(s) &= \frac{L}{2}\cos s, & Y_{\mathrm{I}L}(s) &= \frac{L}{2}\sin s \\ X_{\mathrm{II}}(\sigma,\tau) &= 2L + L\cos\sigma\cos\tau, & Y_{\mathrm{II}}(\sigma,\tau) &= -L\cos\sigma\sin\tau, \\ X_{\mathrm{II}L}(s) &= L + \frac{L}{2}\cos s, & Y_{\mathrm{I}L}(s) &= -\frac{L}{2}\sin s \end{aligned} \tag{15}$$

Strings I and II have equal and opposite angular momenta given by $J_{\mathrm{I}} = L^2/\alpha'$, $J_{\mathrm{II}} = -L^2/\alpha'$. As they rotate, the end $\sigma = 0$ of string I touches the end $\sigma = \pi$ of string II at $\tau = n\pi$, $n =$ integer.

Consider the situation where the strings join at $\tau = 0$. The resulting open string solution has $J = 0$, since angular momentum is conserved and the original total angular momentum of the system is zero. By applying the formulae of [1], we find the solution after joining $X(\tau,\sigma) = X_L(\tau+\sigma) + X_L(\tau-\sigma)$ with

$$X_L(s) = \begin{cases} L + \frac{L}{2}\cos 2s & -\frac{\pi}{2} \leqslant s < \frac{\pi}{2}, \\ -\frac{L}{2}\cos 2s & \frac{\pi}{2} \leqslant s < \frac{3\pi}{2}, \end{cases} \qquad Y_L(s) = -\frac{L}{2}\sin 2s. \tag{16}$$

This solution is shown in figure 1. It describes an open string which at $\tau = 0$ is completely straight, then it bends and contracts until it becomes a point at $\tau = \pi/2$. The solution is periodic with period π.

In the regime that the size of the strings is much larger than the gravitational radius R_s, the evolution is governed by the classical string equations of motion. As the open string reduces its size, gravitational effects become important. A string which reduces to a point should clearly undergo gravitational collapse. This should happen when the size of the string becomes smaller than R_s. For a string of length $\ell \sim M/\mu$, where $\mu = 1/(2\pi\alpha')$ is the string tension, the gravitational radius is $R_s \sim 2G\mu\ell$. Therefore, when the string contracts by a factor of order $(G\mu)^{-1}$, gravitational collapse should be inevitable and a horizon will form.

The starting point of the above example involves two open strings, which we know to be unstable by breaking. The same process can occur for the long-lived strings of section 4, which have an essentially identical four-dimensional dynamics.

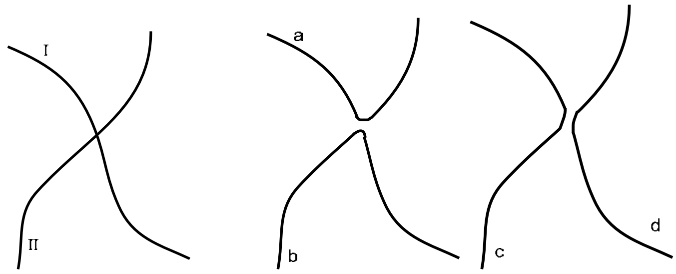

Figure 2. Interconnection process. When two strings cross, there are two possible ways that they can interconnect, leading to strings *a* and *b* or strings *c* and *d*.

5.2. By string interconnection

When two fundamental strings cross, there is a probability given by the string coupling that the strings will interconnect, as in figure 2. This is a common process in 3+1 dimensions, where two infinitely long strings always cross for generic initial data. For finite-size strings, the collision has a cross section of the order of the square of the length of the string.

An interesting question is: what is the probability that a black hole is formed as a result of the collision. Computing this from string perturbation theory is obviously very complicated, so we will try to address this question by means of the following experiment: we send two straight rotating strings against each other, with a random position for the centre-of-mass coordinates and a random value for the relative orientation (within the range where the interconnection is possible). After repeating the experiment N_e times, we ask how many of the resulting string configurations are black holes. We will consider several conditions for black hole formation. One condition is that one of the two final strings completely lie inside its Schwarzschild radius, i.e. $R(\sigma) < R_s$ for all σ, at some time during the evolution. Another condition is that at some time the average size $\bar{R}$ of the string lies inside its Schwarzschild radius. Finally, a third condition is that a segment of the string lies within the Schwarzschild radius. In our study, the reduction to a small size just follows by the natural shrinking of the string that results from flat space evolution, without taking into account the gravity. In any of these three situations, gravitational forces become very strong when the string size approaches R_s and should enhance the evolution towards the collapse.

5.2.1. Black hole events from interconnection. Consider two open strings of (opposite) maximal angular momentum in the XY plane, having the same energy, which cross at some angle at $\tau = 0$. The analytic solutions after the interconnection are obtained by applying the formulae of [1]. We now consider the evolution of each of the two outgoing pieces and study the possible black hole formation. It is convenient to express R_s as $R_s = 2(G\mu)M/\mu$, $\mu = 1/(2\pi\alpha')$. The fundamental string has a tension μ whose value could be anywhere between the TeV scale and the Planck scale. In brane inflation models, one expects a narrower range $10^{-12} < G\mu < 10^{-6}$. The number of black hole events $N_{\rm bh}$ depends on the value of $G\mu$. Table 1 summarizes our results. We see that the condition $R(\sigma) < R_s$ for all σ gives less black hole events. This is due to the cases where a small tail of the string lies outside the Schwarzschild radius. From table 1 one sees that $N_{\rm bh}$ (computed with either criterium) has a power-like dependence with $G\mu$.

A typical black hole event is shown in figure 3. The string after the interconnection has a kink at $\tau = 0$, which then separates into two kinks moving in opposite directions. If the two

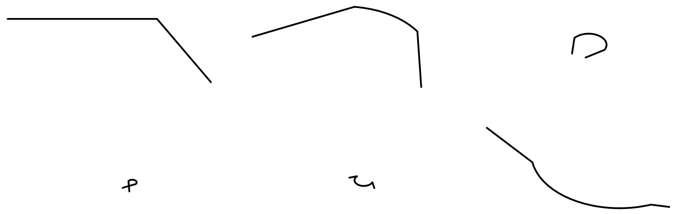

Figure 3. A string analytic solution resulting from interconnection possibly leading to gravitational collapse, after shrinking by its own classical evolution.

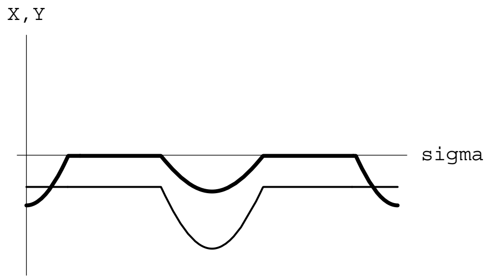

Figure 4. After interconnection, a part of the string shrinks becoming a point during the time evolution. In the figure, these are the values of σ for which both X (thicker line) and Y (thinner) are constant.

Table 1. Number of black hole events in N_e string collisions.

N_e	$G\mu$	N_{bh} $(\bar{R} < R_s)$	N_{bh} $(R(\sigma) < R_s)$
10 000	10^{-2}	1900–2000	1100–1200
10 000	10^{-3}	300–320	95–110
10 000	10^{-4}	40–46	1–3
50 000	10^{-5}	20–30	0–4
50 000	10^{-6}	3–5	0

pieces that form the string have a comparable size, then the strings after the interconnection will have a small angular momentum. This is typically the situation leading to a contraction of the strings to very small size. This figure, however, does not give information on how the mass is distributed. In fact, as we will see in the following section, at some $\tau = \tau_0$ an important fraction of the mass is always concentrated at one point. This fact, which cannot be seen in figure 3 (but can be seen in figure 4), implies that all the cases of this collision of Jmax + antiJmax strings should lead to black hole formation. In particular, this indicates that the cross-section for the scattering of two long strings to form a black hole is essentially given by the geometric area of the overlap of the two strings, times $g_o^4 = g_s^2$, where g_s being the closed string coupling constant.

5.2.2. Inevitable collapse in a generic $J_{\text{I}} = -J_{\text{II}}$ case. In section 5.1, we have already seen that a black hole will form in the case of the *joinings* of Jmax and antiJmax strings. In that case, the string that results from the joining process shrinks, becoming a point at $\tau = \pi/2$.

That is, just by the evolution dynamics in a flat space, all the mass concentrates in a region of the zero size. In the case of interconnection, we find that a *finite* fraction of its arbitrarily large mass shrinks to the zero size at some specific time τ_0. The collapse of a finite mass to a point is a clear sign black hole formation, since gravitational effects can only reinforce the collapse (quantum gravity effects are negligible for large mass black holes, since the horizon has a size much larger than the Planck scale). Therefore, an arbitrarily large black hole is the generic result of the interconnection of arbitrarily large strings of equal and opposite maximal angular momenta. The underlying mechanism is the cancellation of the dependence in σ of the left part with the right part for some value of $\tau = \tau_0$. An example is shown in figure 4.

To conclude, it is possible to follow analytically the string evolution after interconnection, splitting or joining. The evolution exhibits generic features, like kinks. Strings with low angular momenta, like the circular pulsating string or the configurations studied above, tend to contract to a very small size, and appear as potential candidates for black holes.

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