Colliding branes and formation of spacetime singularities in string theory

Andreas Tziolas† and Anzhong Wang‡

GCAP-CASPER, Physics Department, Baylor University, Waco, TX 76798, USA

Zhong Chao Wu∗

Department of Physics, Zhejiang University of Technology, Hangzhou, 310032, China

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Colliding branes without $Z_2$ symmetry and the formation of spacetime singularities in string theory are studied. After developing the general formulas to describe such events, we study a particular class of exact solutions first in the 5-dimensional effective theory, and then lift it to the 10-dimensional spacetime. In general, the 5-dimensional spacetime is singular, due to the mutual focus of the two colliding 3-branes. Non-singular cases also exist, but with the price that both of the colliding branes violate all the three energy conditions, weak, dominant, and strong. After lifted to 10 dimensions, we find that the spacetime remains singular, whenever it is singular in the 5-dimensional effective theory. In the cases where no singularities are formed after the collision, we find that the two 8-branes necessarily violate all the energy conditions.

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I. INTRODUCTION

Branes in string/M-Theory are fundamental constituents [1], and of particular relevance to cosmology [2, 3]. These substances can move freely in the bulk, collide, recoil, reconnect, and whereby, among other possibilities, form a brane gas in the early universe [4], or create an ekpyrotic/cyclic universe [5]. Understanding these processes is fundamental to both string/M-Theory and their applications to cosmology [6].

Recently, Maeda and his collaborators numerically studied the collision of two branes in a five-dimensional bulk, and found that the formation of a spacelike singularity after the collision is generic [7] (See also [8]). This is an important result, as it implies that a low-energy description of colliding branes breaks down at some point, and a complete predictability is lost, without the complete theory of quantum gravity. Similar conclusions were obtained from the studies of two colliding orbifold branes [9]. However, lately it was argued that, from the point of view of the higher dimensional spacetime where the low effective action was derived, these singularities are very mild and can be easily regularised [10].

Lately, we constructed a class of exact solutions with two free parameters to the five-dimensional Einstein field equations, which represents the collision of two timelike 3-branes [11]. We found that, among other things, spacelike singularities generically develop after the collision, due to the mutual focus of the two branes. Non-singular spacetime can be constructed only in the case where both of the two branes violate the energy conditions.

In this paper, we shall systematically study the collision of two timelike 8-branes without $Z_2$ symmetry in the framework of string theory. In particular, in Section II, starting with the Neveu-Schwarz/Neveu-Schwarz (NS-NS) sector in (D+d) dimensions, $M_{D+d} = M_D \times M_d$, we obtain the D-dimensional effective theory in both the string frame and the Einstein frame, by toroidal compactifications. To study the collision of two branes, we add brane actions to the D-dimensional effective action, and then derive the gravitational and matter field equations, including the ones on the two branes. In Section III, we apply these general formulas to the case where $D = 5 = d$ for a large class of spacetimes, and obtain the explicit field equations both outside the two branes and on the two branes. In Section IV, we construct a class of exact solutions in the Einstein frame, in which the potential of the radion field on the two branes take an exponential form, while the matter fields on the two branes are dust fluids. After identifying spacetime singularities both outside and on the branes, we are able to draw the corresponding Penrose diagrams for various cases. In Sections V, we study the local and global properties of these solutions in the 5-dimensional string frame, while in Section VI we first lift the solutions to 10 dimensions, and then study the local and global properties of these 10-dimensional solutions in details. In Section VII, we derive our main conclusions and present some remarks. There is also an Appendix, in which we study a class of 10-dimensional spacetimes. In particular, we divide the Einstein tensor explicitly into three parts, one on each side of a colliding brane, and the other is on the brane. It is remarkable that the part on the brane can be written in the form of an anisotropic fluid.

II. THE MODEL

Let us consider the toroidal compactification of the NS-NS sector of the action in (D+d) dimensions, $M_{D+d} = \ldots$
spacetimes described by the metric, where in this paper we consider the \((D + d)\)-dimensional spacetimes described by the metric,

\[
S_{D+d} = -\frac{1}{2\kappa_*^{D+d}} \int d^{D+d} x \sqrt{|\mathcal{g}_{D+d}|} e^{-\hat{F}} \left\{ \hat{R}_{D+d} [\hat{g}] + \frac{1}{6} \hat{g}^{AB} \left( \nabla_A \hat{\Phi} \right) \left( \nabla_B \hat{\Phi} \right) - \frac{1}{12} \hat{H}^2 \right\},
\]

where in this paper we consider the \((D + d)\)-dimensional spacetimes described by the metric,

\[
d\hat{s}_{D+d}^2 = \hat{g}_{AB} dx^A dx^B = \gamma_{ab} (x^c) dx^a dx^b + \hat{\sigma}^2 (x^c) \hat{\gamma}_{ij} (z^k) dz^i dz^j,
\]

with \(\gamma_{ab} (x^c)\) and \(\hat{\sigma} (x^c)\) depending only on the coordinates \(x^a\) of the spacetime \(M_D\), and \(\hat{\gamma}_{ij} (z^k)\) only on the internal coordinates \(z^k\), where \(a, b, c = 0, 1, 2, \ldots, D - 1; \ i, j, k = D, D + 1, \ldots, D + d - 1\); and \(A, B, C = 0, 1, 2, \ldots, D + d - 1\). Assuming that matter fields are all independent of \(z^k\), one finds that the internal space \(M_d\) must be Ricci flat,

\[
D \equiv \mathcal{R}_{(\gamma)} = 0.
\]

For the purpose of the current work, it is sufficient to assume that \(M_d\) is a \(d\)-dimensional torus, \(T^d = S^1 \times S^1 \times \ldots \times S^1\). Then, we find that

\[
\hat{R}_{D+d} [\hat{g}] = R_D [\gamma] + \frac{d(d-1)}{\hat{\sigma}^2} \gamma^{ab} \left( \nabla_a \hat{\phi} \right) \left( \nabla_b \hat{\phi} \right) - \frac{2}{\hat{\sigma}^2} \gamma^{ab} \left( \nabla_a \nabla_b \hat{\sigma} \right).
\]

Ignoring the dilaton \(\hat{\Phi}\) and the form fields \(\hat{H}\),

\[
\hat{\phi} = 0 = \hat{H},
\]

we find that the integration of the action (2.1) over the internal space yields,

\[
S_D = -\frac{1}{2\kappa_0^D} \int d^D x \sqrt{|\gamma|} \mathcal{G}^{d} \left\{ R_D [\gamma] + \frac{d(d-1)}{\hat{\sigma}^2} \gamma^{ab} \left( \nabla_a \hat{\phi} \right) \left( \nabla_b \hat{\phi} \right) \right\},
\]

where

\[
\kappa_0^D \equiv \frac{\kappa_2^2 D}{V_s},
\]

and \(V_s\) is defined as

\[
V_s = \int \sqrt{|\gamma|} d^d z.
\]

For a string scale compactification, we have \(V_s = \left( 2\pi \sqrt{\alpha'} \right)^d\), where \(2\pi \alpha'\) is the inverse string tension.

After the conformal transformation,

\[
g_{ab} = \hat{\phi}^{2d} \gamma_{ab},
\]

the \(D\)-dimensional effective action of Eq. (2.3) can be cast in the minimally coupled form,

\[
S^{(E)}_D = -\frac{1}{2\kappa_0^D} \int d^D x \sqrt{|g_D|} \left\{ R_D [g] - \kappa_0^2 (\nabla \phi)^2 \right\},
\]

where

\[
\phi \equiv \pm \left( \frac{(D + d - 2)d}{\kappa_0^2 (D - 2)} \right)^{1/2} \ln \left( \hat{\phi} \right).
\]

The action of Eq. (2.7) is usually referred to as the string frame, and the one of Eq. (2.10) as the Einstein frame. It should be noted that solutions related by this conformal transformation can have completely different physical and geometrical properties in the two frames. In particular, in one frame a solution can be singular, while in the other it can be totally free from any kind of singularities. A simple example is the flat FRW universe which is always conformally flat, \(\gamma_{ab} = a^2 (t) \eta_{ab}\). But the spacetime described by \(\gamma_{ab}\) usually has a big bang singularity, while the one described by \(\eta_{ab}\) is Minkowski, and does not have any kind of spacetime singularities.

To study the collision of two branes, we add the following brane actions to \(S^{(E)}_D\) of Eq. (2.10),

\[
S^{(E,I)}_{D-1,m} = \int_{M^{(I)}_{D-1}} \sqrt{|g^{(I)}_{D-1}|} \left( L^{(m,1)}_{D-1}(\psi) - V^{(I)}_{D-1}(\phi) \right) \times d^{D-1} \xi_{(I)},
\]

where \(I = 1, 2, \ V^{(I)}_{D-1}(\phi)\) denotes the potential of the scalar field \(\phi\) on the \(I\)-th brane, and \(\xi_{\mu}^{(I)}\)'s are the intrinsic coordinates of the \(I\)-th brane, where \(\mu, \nu, \lambda = 0, 1, 2, \ldots, D - 2\). \(L^{(m,1)}_{D-1}(\psi)\) is the Lagrangian density of matter fields located on the \(I\)-th brane, denoted collectively by \(\psi\). It should be noted that the above action does not include kinetic terms of the scalar field on the branes. This setup is quite similar to the Horava-Witten heterotic M-Theory on \(S^1/Z_2\) [13, 14], in which the two potentials \(V^{(I)}_{1}(\phi)\) and \(V^{(I)}_{2}(\phi)\) have opposite signs. It is also similar to the modulus stabilization mechanism of Goldberger and Wise [15], which has been lately applied to orbifold branes in string theory [16]. The two branes are localized on the surfaces,

\[
\Phi_I (x^\alpha) = 0,
\]

or equivalently

\[
x^\alpha = x^\alpha \left( \xi_{\mu}^{(I)} \right).
\]

\(V^{(I)}_{D-1}\) denotes the determinant of the reduced metric \(g^{(I)}_{\mu\nu}\) of the \(I\)-th brane, defined as

\[
g^{(I)}_{\mu\nu} \equiv g_{\mu\nu} \xi^{(I)}_{\nu} \xi^{(I)\mu} \xi_{(I)}^{-1},
\]
where

\[ e^{(I)\,a}_{(\mu)} \equiv \frac{\partial x^a}{\partial \xi_{(I)}^{(\mu)}} \mid_{M^{(I)}_{D-1}} \]  \hspace{1cm} (2.16)

Then, the total action is given by,

\[ S^{(E)}_{total} = S^{(E)}_D + \sum_{i=1}^2 S^{(E, I)}_{D-1,m} \]  \hspace{1cm} (2.17)

Then, the non-vanishing components of the Ricci tensor is given by

\[ R_{tt} = - \{ 3\omega_{tt} + \sigma_{tt} + 3\omega_t(\omega_t - \sigma_t) \]  
\[ - \sigma_{yy} - 3\omega_y\sigma_y \}, \]  \hspace{1cm} (3.2)

\[ R_{ty} = - \{ \omega_{ty} + \omega_t\omega_y - \omega_t\sigma_y - \omega_y\sigma_t \}, \]  \hspace{1cm} (3.3)

\[ R_{yy} = - \{ 3\omega_{yy} + \sigma_{yy} + 3\omega_y(\omega_y - \sigma_y) \]  
\[ - \sigma_{tt} - 3\omega_t\sigma_t \}, \]  \hspace{1cm} (3.4)

\[ R_{ij} = \delta_{ij} e^{2(\omega_\sigma)} \{ \omega_{tt} + 3\omega_t^2 \]  
\[ - (\omega_{yy} + 3\omega_y^2) \} \]  \hspace{1cm} (3.5)

where now \( i, j = 2, 3, 4, \) and \( \omega_\sigma \equiv \partial \omega /\partial t, \) etc.

We assume that the two colliding 3-branes move along the hypersurfaces given, respectively, by

\[ \Phi_1(t,y) = t - ay = 0, \]  \hspace{1cm} (3.6)

\[ \Phi_2(t,y) = t + by = 0, \]  \hspace{1cm} (3.7)

in order for the two hypersurfaces to be timelike. The two colliding branes divide the whole spacetime into four regions, \( I - IV, \) which are defined, respectively, as

\[ \text{Region I} \equiv \{ x^a : \Phi_1 < 0, \Phi_2 < 0 \}, \]  \hspace{1cm} (3.8)

\[ \text{Region II} \equiv \{ x^a : \Phi_1 > 0, \Phi_2 < 0 \}, \]  \hspace{1cm} (3.8)

\[ \text{Region III} \equiv \{ x^a : \Phi_1 < 0, \Phi_2 > 0 \}, \]  \hspace{1cm} (3.8)

\[ \text{Region IV} \equiv \{ x^a : \Phi_1 > 0, \Phi_2 > 0 \}, \]  \hspace{1cm} (3.8)

as shown schematically in Fig. 11. In each of these regions, we define

\[ F^A = F(t,y)\mid_{\text{Region} A}, \]  \hspace{1cm} (3.9)

where now \( A = I, II, III, IV. \)

We also define the two hypersurfaces \( \Sigma_1 \) and \( \Sigma_2 \) as,

\[ \Sigma_1 \equiv \{ x^a : \Phi_1 = 0 \}, \]  \hspace{1cm} (3.10)

\[ \Sigma_2 \equiv \{ x^a : \Phi_2 = 0 \}. \]  \hspace{1cm} (3.10)

Then, it can be shown that the normal vectors to each of these two surfaces are given by

\[ n_a = N (\delta_a^0 - a\delta_a^0), \]  \hspace{1cm} (3.11)

\[ l_a = L (\delta_a^0 + b\delta_a^0), \]  \hspace{1cm} (3.11)

where

\[ F^{(t)} = F(t,y)\mid_{\Phi_t = 0}, \]  \hspace{1cm} (3.11)

\[ N = e^{2(\omega(t))}, \]  \hspace{1cm} (3.11)

\[ L = e^{(2b)}, \]  \hspace{1cm} (3.11)

\[ (a^2 - 1)^{1/2}, \]  \hspace{1cm} (3.11)

\[ (b^2 - 1)^{1/2}. \]  \hspace{1cm} (3.11)
with \( F = \{ \sigma, \omega, \phi \} \). We also introduce the two timelike vectors \( u_a \) and \( v_a \) via the relations,
\[
\begin{align*}
u_a &= \mathcal{N} (a \delta_t^a - \delta_y^a), \\
v_a &= L (b \delta_t^a + \delta_a^y).
\end{align*}
\]
(3.13)
It can be shown that these vectors have the following properties,
\[
\begin{align*}
n_a u^a &= -1 = l_a t^a, \\
u_a u^a &= +1 = v_a v^a, \\
n_a v^a &= 0 = l_a v^a.
\end{align*}
\]
(3.14)

In the following, we shall consider field equations, (2.18) and (2.20), in Regions I – IV and along the hypersurfaces \( \Sigma_{1,2} \), separately.

It should be noted that in the above setup, the two 3-branes do not have the \( Z_2 \) symmetry, in contrast to the setup of Horava-Witten in M theory [13] and of Randall-Sundrum [17].

\section*{A. Field Equations in Regions I – IV}

In these regions, the field equations of Eqs. (2.18) and (2.20) take the form,
\[
R^{A}_{ab} = \varphi^{A}_{ab} \varphi_{,b},
\]
(3.15)
\[
\Box^{(A)} \varphi^{A} = 0,
\]
(3.16)
where \( \varphi = \kappa \delta \), and \( \Box^{(A)} \equiv g^{A \, ab} \nabla^{(A)}_{a} \nabla^{(A)}_{b} \), and \( \nabla^{(A)}_{a} \) denotes the covariant derivative with respect to \( g^{A \, ab} \), and \( g^{A \, ab} \) is the metric defined in Region A. From Eq. (3.15) and the fact that \( \varphi = \varphi(t, y) \), we find that
\[
\omega = \frac{1}{3} \ln (f(t + y) + g(t - y)),
\]
(3.17)
where \( f(t + y) \) and \( g(t - y) \) are arbitrary functions of their indicated arguments. Note that in writing Eq. (3.17) we dropped the super indices \( A \). In the following we shall adopt this convention, except for the case where confusions may raise. In the following we consider only the case where
\[
f' g' \neq 0,
\]
(3.18)
where a prime denotes the ordinary derivative with respect to the indicated argument. Then, introducing two new variables \( \xi_{\pm} \) via the relations,
\[
\xi_{\pm}(t, y) \equiv f(t + y) \pm g(t - y),
\]
(3.19)
we find that Eq. (3.15) yields,
\[
M_+ = \frac{1}{2} \xi_{+} \left( \varphi_{+}^2 + \varphi_{-}^2 \right),
\]
(3.20)
\[
M_- = \xi_+ \varphi_+ \varphi_-, \quad (3.21)
\]
and
\[
M_{++} - M_{--} = -\frac{1}{2} \left( \varphi_{+}^2 - \varphi_{-}^2 \right),
\]
(3.22)
where \( M_{\pm} \equiv \partial M / \partial \xi_{\pm} \), and
\[
M(\xi_+, \xi_-) = \sigma + \frac{1}{3} \ln \xi_+ - \frac{1}{2} \ln (4f'g').
\]
(3.23)
On the other hand, Eq. (3.16) can be cast in the form,
\[
\varphi_{++} - \varphi_{--} + \frac{1}{\xi_+} \varphi_+ = 0.
\]
(3.24)
It should be noted that Eqs. (3.20), (3.22) and (3.24) are not all independent. In fact, Eq. (3.22) is the integrability condition of Eqs. (3.20) and (3.24), and can be obtained from Eqs. (3.20), (3.21) and (3.24). Therefore, in Regions I – IV, the field equations reduce to Eqs. (3.20), (3.21) and (3.24).

To find solutions, one may first integrate Eq. (3.24) to find \( \varphi \), and then integrate Eqs. (3.20) and (3.21) to find \( M \). However, Eq. (3.24) has infinite numbers of solutions, and the corresponding general solutions of \( M \) has not been worked out yet [18]. Once \( \varphi \) and \( M \) are known, the metric coefficients \( \sigma \) and \( \omega \) are then given by
\[
\sigma = M - \frac{1}{3} \ln (f + g) + \frac{1}{2} \ln (4f'g'),
\]
(3.25)
\[
\omega = \frac{1}{3} \ln (f + g).
\]
B. Field Equations on the 3-branes

1. Field Equations on the surface \( \Phi_1 = 0 \)

Across the hypersurface \( \Phi_1 = 0 \), for any given \( C^0 \) function \( F(t, y) \), it can be written as \( \Pi \),

\[
F(t, y) = F^+(t, y)H(\Phi_1) + F^-(t, y) [1 - H(\Phi_1)],
\]

(3.26)

where \( F^+ (F^-) \) denotes the function \( F(t, y) \) defined in the region \( \Phi_1 > 0 (\Phi_1 < 0) \), and \( H(x) \) denotes the heaviside function, defined as

\[
H(x) = \begin{cases} 
1, & x > 0, \\
0, & x < 0.
\end{cases}
\]

(3.27)

On the other hand, projecting \( F_\alpha \) onto the \( n_\alpha \) and \( u_\alpha \) directions, we find

\[
F_{\alpha} = F_\alpha u_\alpha - F_n n_\alpha,
\]

(3.28)

where

\[
F_u \equiv u^a F_{\alpha}, \quad F_n \equiv n^a F_{\alpha}.
\]

(3.29)

Since \( [F_u]^- = 0 \) due to the continuity of \( F \) across the branes, from the above expressions we find

\[
[F_{\alpha}]^- = -[F_n]^- n_\alpha,
\]

(3.30)

where

\[
[F_{\alpha}]^- \equiv \lim_{\Phi_1 \to 0^+} F_{\alpha}^+ - \lim_{\Phi_1 \to 0^-} F_{\alpha}^-.
\]

(3.31)

Then, we find that

\[
F_t = F^+_t H(\Phi_1) + F^-_t [1 - H(\Phi_1)],
\]

\[
F_y = F^+_y H(\Phi_1) + F^-_y [1 - H(\Phi_1)],
\]

\[
F_{tt} = F^+_t H(\Phi_1) + F^-_t [1 - H(\Phi_1)] - N [F_n]^+ \delta(\Phi_1),
\]

\[
F_{ty} = F^+_y H(\Phi_1) + F^-_y [1 - H(\Phi_1)] + a N [F_n]^+ \delta(\Phi_1),
\]

\[
F_{yy} = F^+_y H(\Phi_1) + F^-_y [1 - H(\Phi_1)] - a^2 N [F_n]^+ \delta(\Phi_1),
\]

(3.32)

where \( \delta(\Phi_1) \) denotes the Dirac delta function. Then, we find that the Ricci tensor given by Eqs. (3.22)-(3.32) can be cast in the form,

\[
R_{ab} = R^+_{ab} H(\Phi_1) + R^-_{ab} [1 - H(\Phi_1)] + R_{ab}^{lm} \delta(\Phi_1),
\]

(3.33)

where \( R^+_{ab} (R^-_{ab}) \) is the Ricci tensor calculated in the region \( \Phi_1 > 0 (\Phi_1 < 0) \), and \( R_{ab}^{lm} \) denotes the Ricci tensor calculated on the hypersurface \( \Phi_1 = 0 \), which has the following non-vanishing components,

\[
R_{tt}^{lm} = N \left\{ 3 \left[ \omega_n \right]^+ - \left( a^2 - 1 \right) [\sigma_n]^+ \right\},
\]

\[
R_{ty}^{lm} = -3a N \left[ \omega_n \right]^+,
\]

\[
R_{yy}^{lm} = N \left\{ 3a^2 \left[ \omega_n \right]^+ + \left( a^2 - 1 \right) [\sigma_n]^+ \right\},
\]

\[
R_{ij}^{lm} = N e^{2(\omega^{(1)} - \sigma^{(1)})} \left( a^2 - 1 \right) [\omega_n]^+ \delta_{ij}.
\]

(3.34)

On the hypersurface \( \Phi_1 = 0 \), the metric \( \Pi \) reduces to

\[
ds^2_{\Phi_1=0} = g^{(1)}_{\mu\nu} dx^{(1)}_{\mu} dx^{(1)}_{\nu} = d\tau^2 - a_\alpha^2(\tau) d\Sigma_0^2,
\]

(3.35)

where \( \xi^\mu \equiv \{ \tau, x^2, x^3, x^4 \} \), and

\[
d\tau = \epsilon_\tau \left( \frac{a^2 - 1}{a^2} \right)^{1/2} e^{\sigma^{(i)}} dt,
\]

\[
a_\alpha(\tau) = e^{\xi^{(i)}},
\]

(3.36)

with \( \epsilon_\tau = \pm 1 \). Then, we find that

\[
\epsilon^{(1)}_{(\tau)} \equiv \frac{\partial \xi^{(i)}}{\partial \tau} = \epsilon_\tau \left( \frac{\delta^2 + 1}{a^2} \delta_y \right),
\]

\[
\epsilon^{(1)}_{(\tau)} \equiv \frac{\partial \xi^{(i)}}{\partial \tau} = \delta_y,
\]

\[
\sqrt{g_4^{(1)}} = e^{-2\sigma^{(1)}},
\]

(3.37)

where \( i = 2, 3, 4 \) and \( i \equiv dt/d\tau \). Then, the field equations of Eq. (2.20) can be written as

\[
[\omega_n]^+ = \frac{\kappa^2 e^{-\sigma^{(1)}}}{3 (a^2 - 1)^{1/2}} \left( \rho_m^{(1)} + \nabla_4^{(1)} \right),
\]

(3.38)

\[
2 [\omega_n]^+ + [\sigma_n]^+ = \frac{\kappa^2 e^{-\sigma^{(1)}}}{(a^2 - 1)^{1/2}} \left( \nabla_4^{(1)} - p_m^{(1)} \right),
\]

(3.39)

where in writing the above expressions we had assumed that \( T^{(m,1)}_{\mu\nu} \) takes the form of a perfect fluid,

\[
T^{(m,1)}_{\mu\nu} \equiv \left( \rho_m^{(1)} + p_m^{(1)} \right) w_{\mu}^{(1)} w_{\nu}^{(1)} - p_m^{(1)} g_{\mu\nu},
\]

\[
w_{\mu}^{(1)} = \delta_\mu^\gamma.
\]

(3.40)

Similarly, it can be shown that the Klein-Gordon equation (2.20) and the conservation law of the matter fields (2.21) on \( \Sigma_1 \) take, respectively, the forms,

\[
[\phi_n]^+ = \frac{e^{-\sigma^{(1)}}}{(a^2 - 1)^{1/2}} \frac{\partial \nabla_4^{(1)}}{\partial \phi},
\]

\[
\frac{dp_m^{(1)}}{d\tau} + 3H_u \left( \rho_m^{(1)} + p_m^{(1)} \right) = 0,
\]

(3.41)

where \( H_u \equiv \dot{a}_u/a_u \).
2. Field Equations on the Surface \( \Phi_2 = 0 \)

Following a similar procedure as what we did in the last sub-section, one can show that the Ricci tensor across the brane \( \Phi_2 = 0 \) can be written as

\[
R_{ab} = R^+_{ab}H(\Phi_2) + R^-_{ab}[1 - H(\Phi_2)] + R_{ab}^i\delta(\Phi_2),
\]

where \( R^+_{ab} \) (\( R^-_{ab} \)) now is the Ricci tensor calculated in the region \( \Phi_2 > 0 \) (\( \Phi_2 < 0 \)), and \( R_{ab}^i \) denotes the Ricci tensor calculated on the hypersurface \( \Phi_2 = 0 \), which has the following non-vanishing components,

\[
R_{tt}^m = L \left\{ 3 [\omega t]^- - (b^2 - 1) [\sigma t]^- \right\},
\]

\[
R_{ty}^m = 3bL [\omega t]^-,
\]

\[
R_{yy}^m = L \left\{ 3b^2 [\omega t]^- + (b^2 - 1) [\sigma t]^- \right\},
\]

where \( \omega t \equiv \eta \omega_a \) etc. On the hypersurface \( \Phi_2 = 0 \), the metric (3.34) reduces to

\[
ds^2_{1+5} = g_{\mu\nu}^2d\xi^\nu d\xi^\nu = dt^2 - a^2_{\mu}(\eta)d\Sigma_0^2,
\]

where \( \xi_{(2)} = \{ \eta, x^2, x^3, x^4 \} \), and

\[
d\eta \equiv \eta \left( \frac{b^2 - 1}{b^2} \right)^{1/2} e^{\sigma(\eta)} dt,
\]

\[
a_{\nu}(\eta) \equiv e^{\nu(\eta)},
\]

with \( e_\eta = \pm 1 \). Then, we find that

\[
e^{(2)}(\eta) \equiv \frac{\partial w_{\nu}}{\partial \eta} = t^{\nu} \left( \delta^{\nu}_t - \frac{1}{b} \delta^{\nu}_y \right),
\]

\[
e^{(2)}(\eta) \equiv \frac{\partial w_{\nu}}{\partial \xi^{(2)}} = \delta^{\nu}_i,
\]

\[
\sqrt{|g_2^{(2)}} = e^{-2\sigma(\eta)},
\]

where \( t^* \equiv dt/d\eta \). Hence, the field equations of Eq. (2.18) can be written as

\[
[\omega]^- = \frac{k_5^2 e^{-\sigma(\eta)}}{3 (b^2 - 1)^{1/2} (\rho_{m}^{(2)} + V_4^{(2)})}, \quad (3.48)
\]

\[
2[\omega]^- + [\sigma]^- = \frac{k_5^2 e^{-\sigma(\eta)}}{(b^2 - 1)^{1/2}} \left( V_4^{(2)} - p_{m}^{(2)} \right), \quad (3.49)
\]

where in writing the above equations we had assumed that \( T_{(m,2)}^{(m,2)} \) takes the form,

\[
T_{(m,2)}^{(m,2)} = \left( \rho_{m}^{(2)} + p_{m}^{(2)} \right) w_{\mu}^{(2)} w_{\nu}^{(2)} - p_{m}^{(2)} g_{\mu\nu},
\]

\[
w_{\mu}^{(2)} = \delta_{\mu}^\eta.
\]

Similarly, it can be shown that the Klein-Gordon equation (2.20) and the conservation law of the matter fields (2.21) on \( \Sigma_2 \) take, respectively, the forms,

\[
[\phi t]^- = - \frac{e^{-\sigma(\eta)}}{(b^2 - 1)^{1/2}} \frac{dV_4^{(2)}}{d\phi}, \quad (3.51)
\]

\[
d\rho_m^{(2)} d\eta + 3H_v \left( \rho_m^{(2)} + p_m^{(2)} \right) = 0, \quad (3.52)
\]

where \( H_v \equiv \alpha_v^*/\alpha_v \).

IV. PARTICULAR SOLUTIONS FOR COLLIDING TIMELIKE 3-BRANES IN THE EINSTEIN FRAME

Choosing the potentials \( V_4^{(I)}(\phi) \) on the two branes as

\[
V_4^{(I)}(\phi) = V_4^{(I,0)} e^{-\alpha_\phi}, \quad (4.1)
\]

where \( V_4^{(I,0)} \)'s and \( \alpha \) are constants, and that the matter fields on each of the two branes are dust fluids, i.e.,

\[
p_m^{(I)} = 0, \quad (4.2)
\]

we find a class of solutions, which represents the collision of two timelike 3-branes and is given by

\[
\sigma = \left( \chi^2 - \frac{1}{3} \right) \ln (X_0 - X) + \sigma_0, \quad (4.3)
\]

\[
\omega = \frac{1}{3} \ln (X_0 - X) + \omega_0, \quad (4.4)
\]

where \( \chi \equiv \kappa_5/(\sqrt{2}\alpha) \), \( A_0 \), \( \sigma_0 \), \( \omega_0 \) and \( \phi_0 \) are arbitrary constants, and

\[
X = b(t - ay) H(\Phi_1) + a(t + by) H(\Phi_2)
\]

The constants \( a \) and \( b \) are given by

\[
b(a^2 - 1) = \frac{3\kappa_5^{2} V_4^{(1,0)}}{3\chi^2 + 1},
\]

\[
a(b^2 - 1) = -\frac{3\kappa_5^{2} V_4^{(2,0)}}{3\chi^2 + 1}.
\]

When \( \alpha = \pm \infty \), the solutions reduces to the ones studied previously [11]. So, in the rest of this paper we shall consider only the case where \( \alpha \neq \pm \infty \). Without loss of generality, we can always set \( \sigma_0 = \omega_0 = \phi_0 = 0 \), and assume that

\[
X_0 > 0. \quad (4.6)
\]
It can be shown that the field equations, Eqs. (3.15) and (3.16) [or Eqs. (3.20), (3.21) and (3.24)], in Regions $I - IV$ are satisfied identically for the above solutions. To study the singular behavior of the spacetime in each of the four regions, we calculate the Ricci scalar, which in the present case is given by

$$R = \kappa_s^2 g^{ab} \phi_{,a} \phi_{,b} = \frac{\kappa_s^2 B}{\alpha^2 (X_0 - X)^{2(\chi^2 + 2/3)}}. \quad (4.7)$$

where $X$ is given by Eq. (4.4), and

$$B = \begin{cases} 
(a+b)^2, & \text{IV}, \\
-a^2 (b^2 - 1), & \text{III}, \\
-b^2 (a^2 - 1), & \text{II}, \\
0, & \text{I.} 
\end{cases} \quad (4.8)$$

On the 3-brane located on $\Phi_1 = 0$, the reduced metric takes the form,

$$ds_5^{2} |_{\Sigma_1} = d\tau^2 - a_u^2(\tau)d^2\Sigma_0, \quad (4.9)$$

where

$$a_u(\tau) = \begin{cases} 
\frac{[\beta (\tau_s - \tau)]^{\chi^2 + 2}}{X_0^{1/3}}, & \Phi_2 > 0, \\
\frac{1}{X_0^{1/3}}, & \Phi_2 < 0, 
\end{cases} \quad (4.10)$$

with

$$\Phi_2 |_{\Phi=0} = \frac{a + b}{a} t, \quad \beta \equiv \frac{|a(a + b)|}{(a^2 - 1)^{1/2}} \left(\chi^2 + \frac{2}{3}\right), \quad \tau_s \equiv \beta^{-1} X_0^{\chi^2 + \frac{2}{3}} \quad (4.11)$$

Note that in writing the above expressions, we had chosen $\epsilon_s = \text{sign}(a + b)$. From Eqs. (3.38) and (3.39), on the other hand, we find that

$$\rho_m^{(1)} = \frac{1}{X_0 - X(t)} = \begin{cases} 
\frac{[\beta (\tau_s - \tau)]^{\chi^2 + 2}}{X_0^{1/3}}, & \Phi_2 > 0, \\
\frac{1}{X_0^{1/3}}, & \Phi_2 < 0, \end{cases} \quad (4.12)$$

where

$$\rho_m^{(1)} = \frac{b}{\alpha^2 \kappa_s^2} \left(\frac{2}{3} - \chi^2\right), \quad X(t) \equiv (a + b) t H(\Phi_2). \quad (4.13)$$

From Eqs. (4.3) and (4.4) we also find that

$$\phi^{(1)}(\tau) = \begin{cases} 
\frac{1}{3 \alpha^2 X_0^{1/2}} \ln [\beta (\tau_s - \tau)], & \Phi_2 > 0, \\
\frac{1}{3} \ln X_0, & \Phi_2 < 0. \end{cases} \quad (4.14)$$

Similarly, on the 3-brane located on the hypersurface $\Phi_2 = 0$, the reduced metric takes the form,

$$ds^2_{5} |_{\Sigma_2} = d\eta^2 - a_v^2(\eta)d^2\Sigma_0, \quad (4.15)$$

where

$$a_v(\eta) = \begin{cases} 
\ln \left[\frac{\epsilon_\eta (\eta_s - \eta)}{X_0^{1/3}}\right], & \Phi_1 > 0, \\
X_0^{1/3}, & \Phi_1 < 0, \end{cases} \quad (4.16)$$

with $\epsilon_\eta = \text{sign}(a + b)$, and

$$\phi^{(2)}(\eta) = \begin{cases} 
\frac{1}{\alpha^{3 \chi^2 + 2}} \ln \left[\frac{\epsilon_\eta (\eta_s - \eta)}{X_0^{1/3}}\right], & \Phi_1 > 0, \\
\frac{1}{\alpha} \ln X_0, & \Phi_1 < 0, \end{cases} \quad (4.17)$$

The field equations (3.48) and (3.49), on the other hand, yield

$$\rho_m^{(2)} = \frac{1}{X_0 - X(t)} = \begin{cases} 
\frac{1}{\kappa_s^2} \left(\frac{2}{3} - \chi^2\right), & \Phi_1 > 0, \\
X_0^{1/3}, & \Phi_1 < 0. \end{cases} \quad (4.18)$$

where

$$\rho_m^{(2)} \equiv \frac{a}{\kappa_s^2} \left(\frac{2}{3} - \chi^2\right), \quad X(t) \equiv (a + b) t H(\Phi_2). \quad (4.19)$$

It is interesting to note that when $\chi^2 = 2/3$, we have $\rho_m^{(I)} = 0$, $(I = 1, 2)$, and the two 3-branes are supported only by the tensions $V_4^{(I)}(\phi)$, which are non-zero for any finite value of $\alpha$ [Recall the conditions (3.7)]. It is also remarkable to note that the presence of these two dust fluids is not essential to the singularity nature of the spacetime both in the bulk and on the branes. So, in the following we shall study the case with $\chi^2 = 2/3$ together with other cases.

To study the above solutions further, let us consider the following cases separately: (a) $a > 1$, $b > 1$; (b) $a > 1$, $b < -1$; (c) $a < -1$, $b > 1$; and (d) $a < -1$, $b < -1$.

**A.** $a > 1$, $b > 1$

In this case, from Eq. (4.5) we find that

$$V_4^{(1)}(\phi) > 0, \quad V_4^{(2)}(\phi) < 0, \quad (4.20)$$

while Eqs. (4.12) and (4.18) show that

$$\rho_m^{(1)} = \begin{cases} 
\geq 0, & \chi^2 \leq 2/3, \\
< 0, & \chi^2 > 2/3, \end{cases} \quad (4.21)$$

$$\rho_m^{(2)} = \begin{cases} 
\leq 0, & \chi^2 \leq 2/3, \\
< 0, & \chi^2 > 2/3. \end{cases} \quad (4.21)$$
From Eq. (4.7) we can also see that the spacetime is singular along the line $X_0 = (a + b)t$ in Region IV, the line $X_0 = a(t + by)$ in Region III, and the line $X_0 = b(t - ay)$, as shown in Fig. 2.

Before the collision ($t < 0$), the scalar field is constant, $\phi^{(I)} = \phi_1 \equiv (1/\alpha) \ln(X_0)$, but both of the two potentials $V_4^{(1)}(\phi)$ and $V_4^{(2)}(\phi)$ are not zero, so do the dust energy densities $\rho_m^{(i)}$, except for $\chi^2 = 2/3$. In the case $\chi^2 = 2/3$, the dust fluids disappear and the two branes are supported only by tensions, denoted by the two constant potentials $V_4^{(1)}(\phi_1)$ and $V_4^{(2)}(\phi_1)$, which have the opposite signs, and are quite similar to the case of Randall-Sundrum (RS) branes [17], except for that in the RS model the two branes have $Z_2$ symmetry, while here we do not have. Before the collision, the spacetime on the two branes are flat, that is, the matter fields on the 3-brane do not curve the 3-branes. However, it does curve the spacetime outside the 3-branes. This is quite similar to the so-called self-tuning mechanism of brane worlds [20].

After the collision, the two 3-branes focus each other and finally a spacetime singularity is developed at, respectively, $\tau = \tau_s$ and $\eta = \eta_s$. The spacetime on the two branes is homogeneous and isotropic, and is described, respectively, by Eqs. (4.9)-(4.10) and Eqs. (4.11)-(4.12). The corresponding Penrose diagram is given by Fig. 3.

### B. $a > 1, \ b < -1$

In this case, we find that

$$V_4^{(1)}(\phi) < 0, \quad V_4^{(2)}(\phi) < 0.$$
the 3-brane at the point B because the scalar field and the potentials can be seen from Eqs. (4.7), (4.9) and (4.10). This is for BC the line hypersurfaces, Σ − t, y FIG. 4: The five-dimensional spacetime in the (t, y)-plane for a > −b > 1. The two 3-branes are moving along the hypersurfaces, Σ₁ : t − ay = 0 and Σ₂ : t − |b|y = 0. The spacetime is singular along the line AB in Region IV and the line BC in Region III. The spacetime is also singular on the 3-brane at the point B where τ = τₐ. The four regions, I – IV, are defined by Eq. (3.3).

When −b > a > 1, we have

\[ \Phi_1 |_{\Phi_2 = 0} = -\frac{|b| - a}{|b|} t = \begin{cases} > 0, & t > 0, \\ < 0, & t < 0, \end{cases} \]

\[ \Phi_2 |_{\Phi_1 = 0} = -\frac{|b| - a}{a} t = \begin{cases} < 0, & t > 0, \\ > 0, & t < 0, \end{cases} \]

\[ X_0 - X = \begin{cases} X_0 + (|b| - a) t, & \text{IV}, \\ X_0 - a(t - |b|y), & \text{III}, \\ X_0 + |b|(t - ay), & \text{II}, \\ 0, & \text{I}. \end{cases} \] (4.24)

Then, we find that the spacetime is singular along the line \( X_0 = -(|b| - a) t \) in Region IV and the line \( X_0 = (a - |b|) t \) in Region III, as shown in Fig. 6. Unlike the last case, now the 3-brane on Σ₁ starts to expand at the singular point B where \( \tau = \tauₐ \), as shown in Fig. 6 and collides with the one on Σ₂ at the moment \( \tau = 0 \) (t = 0). After the collision, its energy density \( \rho_m^{(1)} \) the scalar field \( \phi^{(2)} \) and the dust energy density \( \rho_m^{(2)} \) on Σ₂ become time-dependent, and the corresponding spacetime is described by Eqs. (4.15) and (4.19) with \( \eta \in (0, -\infty) \). The corresponding Penrose diagram is given by Fig. 7.
is singular along the line $BC$ in Region $III$. The spacetime is also singular on the 3-brane at the point $B$.

**FIG. 7:** The Penrose diagram for $-b > a > 1$. The spacetime is singular along the lines $AB$ and $BC$.

C. $a < -1, b > 1$

In this case, we find that

$$V_4^{(I)}(\phi) > 0,$$

$$\rho_m^{(I)} = \begin{cases}
  \geq 0, & \chi^2 \leq 2/3, \\
  < 0, & \chi^2 > 2/3,
\end{cases}$$

(4.25)

where $I = 1, 2$. Thus, in contrast to the last case, now both potentials $V_4^{(I)}(\phi)$ are positive, while the two dust energy densities always have the same sign.

1. $-a > b > 1$

When $-a > b > 1$, we have

$$\Phi_1|_{\phi_2=0} = -\frac{|a|-b}{b}t = \begin{cases}
  < 0, & t > 0, \\
  > 0, & t < 0,
\end{cases}$$

$$\Phi_2|_{\phi_1=0} = \frac{|a|}{|\eta|}t = \begin{cases}
  < 0, & t > 0, \\
  > 0, & t < 0,
\end{cases}$$

$$X_0 - X = \begin{cases}
  X_0 + (|a|-b)t, & IV, \\
  X_0 + |a|(t + |a|y), & III, \\
  X_0 - b(t + |a|y), & II, \\
  0, & I.
\end{cases}$$

(4.26)

Then, the spacetime is singular along the line $X_0 = -(|a|-b)t$ in Region $IV$, and along the line $X_0 = b(t + |a|y)$ in Region $II$, as shown in Fig. 9. The corresponding Penrose diagram is given by Fig. 9.

In this case, we also have

$$\phi^{(1)}(\tau) = \begin{cases}
  \frac{1}{\alpha(3\tau^2+2)}\ln[\beta(\tau_0 - \tau)], & t > 0, \\
  \frac{1}{\beta \ln X_0}, & t < 0,
\end{cases}$$

$$\phi^{(2)}(\eta) = \begin{cases}
  \frac{1}{\alpha(3\tau^2+2)}\ln[\gamma(\eta_0 - \eta)], & t > 0, \\
  \frac{1}{\beta \ln X_0}, & t < 0,
\end{cases}$$

$$\rho_m^{(1)} = \begin{cases}
  [\beta(\tau_0 - \tau)]^{-\frac{3}{3\tau^2+2}}, & t > 0, \\
  X_0^{-1}, & t < 0,
\end{cases}$$

$$\rho_m^{(2)} = \begin{cases}
  X_0^{-1}, & t > 0, \\
  [\gamma(\eta_0 - \eta)]^{-\frac{3}{3\tau^2+2}}, & t < 0.
\end{cases}$$

(4.27)

2. $b > -a > 1$

When $b > -a > 1$, we have

$$\Phi_1|_{\phi_2=0} = \frac{b - |a|}{b}t = \begin{cases}
  > 0, & t > 0, \\
  < 0, & t < 0,
\end{cases}$$

$$\Phi_2|_{\phi_1=0} = -\frac{b - |a|}{|a|}t = \begin{cases}
  < 0, & t > 0, \\
  > 0, & t < 0,
\end{cases}$$

$$X_0 - X = \begin{cases}
  X_0 - (b - |a|)t, & IV, \\
  X_0 + |a|(t + |a|y), & III, \\
  X_0 - b(t + |a|y), & II, \\
  0, & I.
\end{cases}$$

(4.28)
We also have

\[
\phi^{(1)}(\tau) = \begin{cases} 
\frac{1}{\alpha} \ln X_0, & t > 0, \\
\frac{1}{\alpha(3\chi^2 + 2)} \ln [\beta (\tau_s - \tau)], & t < 0,
\end{cases}
\]

\[
\phi^{(2)}(\eta) = \begin{cases} 
\frac{1}{\alpha(3\chi^2 + 2)} \ln [\gamma (\eta_s - \eta)], & t > 0, \\
\frac{1}{\alpha} \ln X_0, & t < 0,
\end{cases}
\]

\[
\rho^{(1)}_m = \begin{cases} 
X_0^{-1}, & t > 0, \\
[\beta (\tau_s - \tau)]^{-\frac{1}{3\chi^2 + 2}}, & t < 0,
\end{cases}
\]

\[
\rho^{(2)}_m = \begin{cases} 
[\gamma (\eta_s - \eta)]^{-\frac{1}{3\chi^2 + 2}} X_0^{-1}, & t > 0, \\
X_0^{-1}, & t < 0.
\end{cases}
\]

Then, the spacetime is singular along the line \(X_0 = (b - |a|)t\) in Region IV, and along the line \(X_0 = b(t + |a|y)\) in Region II, as shown in Fig. 10. The corresponding Penrose diagram is given by Fig. 11.

**D. \( a < -1, b < -1 \)**

In this case, we have

\[
V_4^{(1)}(\phi) < 0, \quad V_4^{(2)}(\phi) > 0,
\]

\[
\rho^{(1)}_m = \begin{cases} 
\geq 0, & \chi^2 \geq 2/3, \\
< 0, & \chi^2 < 2/3,
\end{cases}
\]

\[
\rho^{(2)}_m = \begin{cases} 
\geq 0, & \chi^2 \leq 2/3, \\
< 0, & \chi^2 > 2/3.
\end{cases}
\]
Then, we find that branes are supported only by the tensions, where the fluid on each of the two 3-branes disappears, and the Penrose diagram is given by Fig. 13.

The two 3-branes are moving along the hypersurfaces, $\Sigma_1 : t + |a|y = 0$ and $\Sigma_2 : t - |b|y = 0$. The spacetime is free of any kind of spacetime singularities in the four regions, $I - IV$, as well as on the two 3-branes.

It is interesting to note that when $\chi^2 \neq 2/3$, both dust fluid are present, but they always have opposite signs, that is, if one satisfies the energy conditions [21], the other one must violate these conditions.

V. COLLIDING 3-BRANES IN THE 5-DIMENSIONAL STRING FRAME

The spacetime singularity behavior in general can be quite different in the two frames, due to the conformal transformations of Eq.(2.9), which are often singular.

The 5-dimensional spacetime in the string frame is given in Eq.(3.1), and

$\Phi_1 = \frac{|a| + |b|}{|b|} t = \begin{cases} 0, & t > 0, t > 0, \\ -\frac{|a|}{|a|} t, & t > 0, t < 0, \\ \frac{|b|}{|b|} t, & t < 0, t < 0, \end{cases}$

$\Phi_2 = \frac{|a| + |b|}{|a|} t = \begin{cases} 0, & t > 0, t > 0, \\ -\frac{|b|}{|b|} t, & t > 0, t < 0, \\ \frac{|a|}{|a|} t, & t < 0, t < 0, \end{cases}$

$X_0 - X = \begin{cases} X_0 + (|a| + |b|)t, & IV, \\ X_0 + |a| (t - |b|y), & III, \\ X_0 + |b| (t + |a|y), & II, \\ 0, & I. \end{cases}$

Then, we find that

$\phi^{(1)}(\tau) = \begin{cases} \frac{1}{\alpha(3\tau + 2)} \ln [\beta (\tau_s - \tau)], & t > 0, \\ \frac{1}{\alpha} \ln X_0, & t < 0, \end{cases}$

$\phi^{(2)}(\eta) = \begin{cases} \frac{1}{\alpha(3\eta + 2)} \ln [\gamma (\eta_s - \eta)], & t > 0, \\ \frac{1}{\alpha} \ln X_0, & t < 0, \end{cases}$

$\rho^{(1)}_m = \begin{cases} [\beta (\tau_s - \tau)]^{-\frac{3}{3\tau + 2}}, & t > 0, \\ X_0^{-\frac{3}{\alpha}}, & t < 0, \end{cases}$

$\rho^{(2)}_m = \begin{cases} [\gamma (\eta_s - \eta)]^{-\frac{3}{3\eta + 2}}, & t > 0, \\ X_0^{-\frac{3}{\alpha}}, & t < 0. \end{cases}$

Note that in the present case, after the collision $t > 0$, we have $\tau, \eta < 0$. Thus, in this case the spacetime is free of any kind singularity in all the four regions, as well as on the two branes, as shown in Fig. 14. The corresponding Penrose diagram is given by Fig. 15.

It is interesting to note that when $\chi^2 = 2/3$, the dust fluid on each of the two 3-branes disappears, and the branes are supported only by the tensions, where the brane along $\Sigma_1$ has a negative tension, while the one along $\Sigma_1$ has a positive tension. The spacetime is free of any kind singularity in all the four regions, as well as on the two 3-branes.

The 5-dimensional spacetime in the string frame is given by

$$d^2s_5 = \gamma_{ab}dx^a dx^b - e^{2\tilde{\sigma}(t,y)} (dt^2 - dy^2) - e^{2\tilde{\omega}(t,y)} d\Sigma_0^2,$$

where $d\Sigma_0^2$ is given in Eq.(3.1), and

$$\tilde{\sigma}(t,y) = \left(\frac{\chi^2 - \epsilon}{\sqrt{\frac{5}{12}} \chi - \frac{1}{3}}\right) \ln (X_0 - X),$$

$$\tilde{\omega}(t,y) = \left(\frac{1}{\sqrt{\frac{5}{12}} \chi} \right) \ln (X_0 - X),$$

$$\tilde{\phi}(t,y) = \left(X_0 - X\right) \frac{d}{dx} \chi,$$
frame is stronger, and if $\epsilon_\alpha < 0$, we find that the spacetime in the Einstein frame is singular in the string frame whenever it is singular in the Einstein frame, although the strength of the singularity is different, as can be seen clearly from the following expression,

$$\dot{\phi}_a \dot{\phi}^a = \frac{3\chi^2 B}{20 (X_0 - X)^{\frac{\chi}{2} + (\sqrt{\frac{\chi}{6}} - \sqrt{2} \chi)^2}},$$  \hspace{1cm} (5.3)$$

where $B$ is given by Eq. (4.3). Comparing the above expression with Eq. (4.7), we find that the spacetime in Regions I – IV is singular in the string frame whenever it is singular in the Einstein frame, although the strength of the singularity is different, as can be seen clearly from the following expression,

$$20 \frac{\dot{\phi}_a \dot{\phi}^a}{R} = (X_0 - X)^{\chi \sqrt{\frac{\chi}{6}}},$$  \hspace{1cm} (5.4)$$

In particular, if $\epsilon_\alpha > 0$ the singularity in the Einstein frame is stronger, and if $\epsilon_\alpha < 0$ it is the other way around.

2. The Spacetime on the 3-brane $t = ay$

On the hypersurface $t = ay$, the metric (5.1) reduces to

$$d^2 s_5 \big|_{t=ay} = d\tilde{\tau}^2 - a_u^2 (\tilde{\tau}) d\Sigma_0^2,$$  \hspace{1cm} (5.5)$$

where

$$a_u (\tilde{\tau}) = \begin{cases} a_0 (\tilde{\tau}_0 - \tilde{\tau})^{\Delta}, & \Phi_2 > 0, \\ a_0 \tilde{\tau}_0^{\Delta}, & \Phi_2 < 0, \end{cases}$$

with

$$X_0 - X^{(1)} = \begin{cases} \left[ \beta (\tilde{\tau}_0 - \tilde{\tau}) \right]^{\frac{1}{\sqrt{48}}}, & \Phi_2 > 0, \\ \Phi_2 < 0, \end{cases} \hspace{1cm} (5.6)$$

$$\Phi_2|_{\Phi_1=0} = \frac{a + b}{a} t, \hspace{1cm} \beta \equiv \frac{|a(a + b)|}{\sqrt{b^2 - 1}} \delta,$$

$$\hat{\eta}_s \equiv \beta^{-1} X_0^{\frac{3}{2}},$$

$$\Delta \equiv \frac{1}{\delta} \left( \frac{1}{3} - \epsilon \chi \sqrt{\frac{5}{12}} \right).$$  \hspace{1cm} (5.7)$$

Note that in writing the above expressions, we had chosen $\epsilon_\tau = \text{sign}(a + b)$. To study the spacetime singularity on the brane, we calculate the Ricci scalar, which now is given by

$$R^{(4)\lambda}_{\lambda} = \frac{3\Delta (2 - \Delta)}{2a_0 (\tilde{\tau}_0 - \tilde{\tau})^{\Delta + 2}},$$  \hspace{1cm} (5.8)$$

where

$$\Delta + 2 = \frac{1}{\delta} \left[ 2 \left( \epsilon \chi - \sqrt{\frac{15}{64}} \right)^2 + \frac{115}{96} \right] > 0,$$

$$\Delta - 2 = -\frac{1}{\delta} \left[ 2 \left( \epsilon \chi - \sqrt{\frac{5}{48}} \right)^2 + \frac{19}{24} \right] < 0.$$  \hspace{1cm} (5.9)$$

3. The Spacetime on the 3-brane $t = -by$

Similarly, on the 3-brane located on the hypersurface $\Phi_2 = 0$, the metric (5.1) reduces to

$$d^2 s_5 \big|_{t=-by} = d\bar{\eta}^2 - a_v^2 (\bar{\eta}) d\Sigma_0^2,$$  \hspace{1cm} (5.10)$$

where where

$$a_v (\bar{\eta}) = \begin{cases} a_0 (\bar{\eta}_0 - \bar{\eta})^\Delta, & \Phi_1 > 0, \\ a_0 \bar{\eta}_0^\Delta, & \Phi_1 < 0, \end{cases}$$

with

$$X_0 - X^{(2)} = \begin{cases} \left[ \hat{\eta} (\hat{\eta}_0 - \hat{\eta}) \right]^{\frac{1}{\sqrt{48}}}, & \Phi_1 > 0, \\ \Phi_1 < 0, \end{cases} \hspace{1cm} (5.11)$$

$$\Phi_1|_{\Phi_2=0} = \frac{a + b}{b} t, \hspace{1cm} \hat{\eta} \equiv \frac{|a(a + b)|}{\sqrt{b^2 - 1}} \delta,$$

$$\hat{\eta}_s \equiv \hat{\eta}^{-1} X_0^{\frac{3}{2}},$$  \hspace{1cm} (5.12)$$

FIG. 13: The Penrose diagram for $a < -1$, $b < -1$. The spacetime is non-singular in all the regions.
but now we have \( a_0 \equiv \hat{\gamma}^A \) and \( \epsilon_0 = \text{sign}(a + b) \). For the metric (5.10), we also find that

\[
R^{(4)}_{\lambda} = \frac{3\Delta (2 - \Delta)}{2a_0 (\hat{n}_s - \hat{n})^{\Delta + 2}}. \tag{5.13}
\]

From Eqs. (5.8) and (5.13) we can see that the spacetime on each of the branes is not singular when \( \Delta = 0 \) or \( \chi = \sqrt{\frac{4}{15}} \). As a matter of fact, in this case the spacetime on each of the two branes is flat. Thus, in the following we need to consider only the case \( \chi \neq \sqrt{\frac{4}{15}} \).

From Eqs. (5.8) and Eqs. (5.11)-(5.13), it can be shown that the spacetime singularities on each of the two branes are similar to these in the Einstein frame. For example, for the case \( a > 1, b > 1 \), it is singular at \( \tilde{\tau} = \hat{\tau} \), and \( \hat{n} = \hat{n}_s \), which correspond to, respectively, the point \( A \) and \( B \) in Fig. 3. Similarly, the spacetime is free from any kind of singularities for the case \( a < -1, b < -1 \), and the corresponding Penrose diagram is also given by Fig. 3.

**VI. COLLIDING 3-BRANES IN THE 10-DIMENSIONAL SPACETIMES**

Lifting the metric to 10-dimensions, it is given by Eq. (6.1), which can be cast in the form,

\[
d^2 \hat{s}_{10} = \gamma_{ab} dx^a dx^b + \hat{\gamma}^2 (x^c) \hat{\gamma}_{ij} (z^k) dz^i dz^j = e^{2\hat{\sigma}(t,y)} (dt^2 - dy^2) - e^{2\hat{\omega}(t,y)} d\Sigma^2 - \hat{\phi}^2 (t,y) d\Sigma_z. \tag{6.1}
\]

where \( \hat{\sigma}, \hat{\omega} \) and \( \hat{\phi} \) are given by Eq. (5.2), and \( d\Sigma_z^2 = -\sum_{k=1}^{5} \hat{\gamma}_{ij} (z^k) dz^i dz^j \). Then, it can be shown that the spacetime in Regions I - IV is vacuum,

\[
R^{(4)}_{\lambda} = 0, \tag{6.2}
\]

where \( A = I, ..., IV \), as it is expected. To study the singular behavior of the spacetime in these regions, we calculate the Kretschmann scalar, which in the present case is given by

\[
I_{10} = R^{(4)}_{\lambda} R^{4}_{\lambda} = \frac{B^2 I_{10}^{(0)}}{(X_0 - X) (2\chi - \sqrt{\chi^2 + 4})}. \tag{6.3}
\]

where \( B \) is given by Eq. (1.3), and

\[
I_{10}^{(0)} = \frac{1}{45} \left( (720\chi^6 + 1287\chi^4 + 200\chi^2 + 40) - 312\sqrt{\frac{5}{3}} \chi^3 (2 + 3\chi^2) \right). \tag{6.4}
\]

It can be shown that \( I_{10}^{(0)} \) is non-zero for any given \( \chi \). Then, comparing the expression of Eq. (6.3) with Eq. (1.7), we find that the lifted 10-dimensional spacetime has a similar singular behavior as that in the 5-dimensional spacetime in the Einstein frame. In particular, it is also singular on the hypersurface \( X_0 - X = 0 \).

On the hypersurface \( t = ay \), the metric (6.1) reduces to

\[
d^2 \hat{s}_5|_{t=ay} = d\hat{\tau}^2 - a^2 (\hat{\tau}) d\Sigma^2_0 - b^2 (\hat{\tau}) d\Sigma_z^2, \tag{6.5}
\]

where \( a_u (\hat{\tau}) \) and \( b_u (\hat{\tau}) \) are given by Eqs. (5.6) and (5.7). On the 8-brane, the Einstein tensor has distribution given by Eqs. (A.8) and (A.9). Inserting Eq. (5.2) into Eq. (5.21), and noticing that \( \hat{\psi} \equiv \log (\hat{\phi}) \), we find

\[
\hat{\rho}_u = \frac{b (a^2 - 1)}{[X_0 - X^{(1)}(t)]^2}, \tag{6.6}
\]

\[
\hat{\rho}_v = \frac{b (a^2 - 1)}{[X_0 - X^{(1)}(t)]^2} \left( \chi - \epsilon \sqrt{\frac{4}{15}} \right)^2 + \frac{2}{5}, \tag{6.7}
\]

where \( \chi^{(1)}(t) \) is given by Eq. (5.7), and

\[
\mu = 2 \left( \chi - \epsilon \sqrt{\frac{5}{48}} \right)^2 + \frac{1}{8}. \tag{6.7}
\]

Clearly, whenever \( X_0 - X^{(1)}(t) = 0 \), the spacetime on the 8-brane is singular. On the hypersurface \( t = -by \), the metric (6.1) reduces to

\[
d^2 \hat{s}_5 |_{t=-by} = d\hat{\tau}^2 - a_v (\hat{\eta}) d\Sigma^2_0 - b_v (\hat{\eta}) d\Sigma_z^2, \tag{6.8}
\]

where \( a_v (\hat{\eta}) \) and \( b_v (\hat{\eta}) \) are given by Eqs. (5.11) and (5.12). On this 8-brane, the Einstein tensor has distribution given by Eqs. (A.11) and (A.12), which in the present case yield,

\[
\hat{\rho}_v = \frac{a (b^2 - 1)}{[X_0 - X^{(2)}(t)]^2}, \tag{6.9}
\]

\[
\hat{\rho}_v = \frac{a (b^2 - 1)}{[X_0 - X^{(2)}(t)]^2} \left( \chi - \epsilon \sqrt{\frac{4}{15}} \right)^2 + \frac{2}{5}, \tag{6.9}
\]

where \( X^{(2)}(t) \) is given by Eq. (5.12). Thus, the spacetime on this 8-brane is also singular whenever \( X_0 - X^{(2)}(t) = 0 \).

When \( a > 1 \) and \( b > 1 \), from Eqs. (6.6) and (6.8), it can be shown that both of the weak and dominant energy conditions [21] are satisfied by the matter fields on the two 8-branes, provided that

\[
\begin{align*}
\sqrt{\frac{1}{2}} - \sqrt{\frac{1}{3}} & \leq \chi \leq \sqrt{\frac{3}{4}}, & \epsilon = +1, \\
-\sqrt{\frac{3}{2}} & \leq \chi \leq \sqrt{\frac{3}{2}} - \sqrt{\frac{1}{15}}, & \epsilon = -1.
\end{align*} \tag{6.10}
\]
but the strong energy condition is always violated. When 
\( a > 1 \) and \( b < -1 \), the matter field on the 8-brane \( \Phi_1 = 0 \) 
violates all the three energy conditions, while the one 
on the 8-brane \( \Phi_2 = 0 \) satisfies the weak and dominant 
energy conditions, provided that the conditions \( \text{WEC} \) 
holds, but violates the strong one. When \( a < -1 \) and 
\( b > 1 \), it is the other way around, that is, the matter 
field on the 8-brane \( \Phi_1 = 0 \) satisfies the weak and dominant 
energy conditions, provided that the conditions \( \text{SCE} \) 
holds, but violates the strong one, while the one on the 
8-brane \( \Phi_2 = 0 \) violates all the three energy conditions. 
When \( a < -1 \) and \( b < -1 \), the matter fields on the two 8-branes 
all violate the three energy conditions. However, 
in all these four cases, the spacetime singular behavior is 
similar to the corresponding 5-dimensional cases in the 
Einstein frame. In particular, in the first three cases 
the spacetime in the four regions and on the 8-branes 
are always singular, and the corresponding Penrose 
diagrams are given, respectively, by Figs. \([3][5] [7][9] \) and \([11]\) 
but now each point in these figures now represents a 8-
dimensional spatial space. In the last case, in which the 
matter fields on the two 8-branes violate all the energy 
conditions, the spacetime is free of any kind of spacetime 
singularities, either in Regions I – IV or on the two 8-
branes, and the corresponding Penrose diagram is given 
by Fig. \([13]\). Therefore, all the above results seemingly 
indicate that violating the energy conditions is a neces-
sary condition for spacetimes of colliding branes to be 
non-singular.

**VII. CONCLUSIONS**

In this paper, we have first developed the general for-
mulas to describe the collision of two timelike (D-1)- 
branes without \( Z_2 \) symmetry in a D-dimensional 
effective theory, obtained from the toroidal compactification 
of the Neveu-Schwarz/Neveu-Schwarz (NS-NS) sector in 
(D+4) dimensions. Applying the formulas to the case 
\( D = 5 = d \) for a class of spacetimes. In Section III we 
have obtained explicitly the field equations both outside 
and on the 3-branes in terms of distributions. In Section 
IV, we have considered a class of exact solutions that 
represents the collision of two 3-branes in the Einstein frame, 
and studied their local and global properties in details. We 
have found, among other things, that the collision in 
general ends up with the formation of spacetime singular-
ities, due to the mutual focus of the colliding branes, 
although non-singular spacetime also exist, with the price 
that both of the two branes violate all the energy condi-
tions, weak, strong and dominant. Similar conclusions 
hold also in the 5-dimensional string frame. This has 
been done in Section V. In Section VI, after lifted the 
solutions to 10-dimensional spacetimes, we have found 
that the corresponding solutions represent the collision 
of two timelike 8-branes without \( Z_2 \) symmetry. In some 
cases the two 8-branes satisfy the weak and dominant 
energy conditions, while in other case, they do not. But, 
in all these cases the strong energy condition is always 
violated. The formation of spacetime singularities due 
to the mutual focus of the two colliding branes occurs 
general, although the non-singular cases also exist with 
the price that both of the two singular cases also exist 
all the three energy conditions. The spacetime singular behavior is 
similar in the 5-dimensional effective theory to that of 
10-dimensional string theory.

In this paper, we have ignored the dilaton \( \Phi \) and the 
three-form field \( H_{ABC} \). It would be very interesting to 
see how these fields affect the formation of the spacetime 
singularities. In addition, it would also be very interest-
ting to see what might happen if the branes are allowed 
to collide more than one time.

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**Appendix: Gravitational field equations in the 
10-dimensional bulk and on the 8-branes**

For the metric,

\[
d^2 s_{10} = e^{2\sigma(t,y)} (dt^2 - dy^2) - e^{2\omega(t,y)} d\Sigma_0^2 \\
- \dot{\phi}^2 (t,y) d\Sigma_5^2, 
\]

(A.1)

where

\[
d\Sigma_0^2 = \sum_{p=2}^4 (dx^p)^2, \quad d\Sigma_5^2 = \sum_{i=1}^5 (dz^i)^2, 
\]

(A.2)

the non-vanishing components of the Einstein tensor are 
given by,

\[
G_{tt}^{(10)} = 3\dot{\omega}_t (\dot{\sigma}_t + \dot{\omega}_t) + 5\psi_t (\dot{\sigma}_t + 3\dot{\omega}_t + 2\dot{\psi}_t) \\
- 3\dot{\omega}_y y - 5\dot{\psi}_y y - 15\dot{\psi}_t y (\dot{\omega}_y + \dot{\psi}_y) \\
+ \sigma_y (3\dot{\omega}_y + 5\dot{\psi}_y) - 6\dot{\omega}_y^2,
\]

\[
G_{ty}^{(10)} = -3\dot{\omega}_ty - 5\dot{\psi}_ty \\
+ 3 (\dot{\sigma}_t \dot{\omega}_y + \dot{\sigma}_y \dot{\omega}_t - \omega_y \dot{\omega}_t) \\
+ 5 (\dot{\sigma}_t \dot{\psi}_y + \dot{\sigma}_y \dot{\psi}_t - \dot{\psi}_t \dot{\psi}_y),
\]

\[
G_{yy}^{(10)} = -3\dot{\omega}_y y - 5\dot{\psi}_y y - 15\dot{\psi}_t y (\dot{\omega}_y + \dot{\psi}_y) \\
+ \sigma_t (3\dot{\omega}_t + 5\dot{\psi}_t) - 6\dot{\omega}_t^2 \\
+ 3\dot{\omega}_y (\dot{\sigma}_y + \dot{\omega}_y) \\
+ 5\dot{\psi}_y (\dot{\sigma}_y + 3\dot{\psi}_y + 2\dot{\psi}_y),
\]

\[
G_{pq}^{(10)} = \delta_{pq} e^{2(\dot{\omega} - \dot{\sigma})} \left[ \dot{\sigma}_{yy} + 2\dot{\omega}_{yy} + 5\dot{\psi}_{yy} \\
+ 5\dot{\psi}_y (2\dot{\omega}_y + 3\dot{\psi}_y) + 3\dot{\omega}_y^2 \right].
\]
we find that Eq. (A.6) can be cast in the form,
\[
G_{ab}^{(10)Im} = \kappa_2^{10} \left( \rho_u \hat{u}_a \hat{u}_b + \rho^X_u \sum_{p=2}^{4} X_a^{(p)} X_b^{(p)} \right) + \rho_{\hat{z}} \sum_{i=1}^{5} Z_a^{(i)} Z_b^{(i)} ,
\]
where
\[
\rho_u = \frac{1}{N \kappa_2^{10}} \left( 3 [\hat{\omega}_n]^{-5} + 5 \left[ \hat{\psi}_n \right]^{-5} \right) ,
\]
\[
\rho^X_u = \frac{1}{N \kappa_2^{10}} \left( [\hat{\sigma}_n]^{-5} + 2 [\hat{\omega}_n]^{-5} + \left[ \hat{\psi}_n \right]^{-5} \right) ,
\]
\[
\rho_{\hat{z}} = \frac{1}{N \kappa_2^{10}} \left( [\hat{\sigma}_n]^{-5} + 3 [\hat{\omega}_n]^{-5} + 4 \left[ \hat{\psi}_n \right]^{-5} \right) .
\]

**B. Field Equations on the hypersurface \( \Phi_2 = 0 \)**

Similarly, it can be shown that, crossing the hypersurface \( \Phi_2 = 0 \) \( (\Phi_2 < 0) \), and \( G^{(10)Im} \) denotes the distribution of the Einstein tensor on the hypersurface \( \Phi_2 = 0 \), which can be written in the form,
\[
G_{ab}^{(10)Im} = \kappa_2^{10} \left( \rho_v \hat{v}_a \hat{v}_b + \rho^X_v \sum_{p=2}^{4} X_a^{(p)} X_b^{(p)} \right) + \rho_{\hat{z}} \sum_{i=1}^{5} Z_a^{(i)} Z_b^{(i)} ,
\]
where
\[
\rho_v = \frac{1}{L \kappa_2^{10}} \left( 3 [\hat{\omega}_n]^{-5} + 5 \left[ \hat{\psi}_l \right]^{-5} \right) ,
\]
\[
\rho^X_v = \frac{1}{L \kappa_2^{10}} \left( [\hat{\sigma}_l]^{-5} + 2 [\hat{\omega}_n]^{-5} + 5 \left[ \hat{\psi}_l \right]^{-5} \right) ,
\]
\[
\rho_{\hat{z}} = \frac{1}{L \kappa_2^{10}} \left( [\hat{\sigma}_l]^{-5} + 3 [\hat{\omega}_n]^{-5} + 4 \left[ \hat{\psi}_l \right]^{-5} \right) .
\]

and
\[
X_a^{(p)} = e^{\hat{\omega}_n} \hat{\delta}_a^{(2)} , \quad Z_a^{(i)} = e^{\hat{\omega}_n} \hat{\delta}_a^{(i)} ,
\]
\[
\hat{l}_a = \hat{L} (\hat{\delta}_a^{(2)} + \hat{b} \delta_a^{(2)}) , \quad \hat{v}_a = \hat{L} (\hat{b} \delta_a^{(2)} + \delta_a^{(2)}) ,
\]
\[
\hat{L} = \frac{e^{\hat{\omega}_n}}{(b^2 - 1)^{1/2}} .
\]
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