Induced gravity on intersecting brane-worlds
Part II: Cosmology

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Abstract

We explore cosmology of intersecting braneworlds with induced gravity on the branes. We find the cosmological equations that control the evolution of a moving codimension one brane and a codimension two brane that sits at the intersection. We study the Friedmann equation at the intersection, finding new contributions from the six dimensional bulk. These higher dimensional contributions allow us to find new examples of self-accelerating configurations for the codimension two brane at the intersection and we discuss their features.
1 Introduction

Brane-world models offer new perspectives for explaining the present day acceleration in purely geometrical terms, without the need to introduce dark energy [1,2,3] (for a review see [4]). A celebrated example is the Dvali-Gabadadze-Porrati (DGP) model in a 5d spacetime [1]. The brane action includes a quantum-induced Einstein-Hilbert action that recovers 4d gravity on small scales. This model realizes a so-called self-accelerating solution that features a 4d de Sitter phase even though the 3-brane is completely empty. However, so far, only codimension-one examples of such solutions have been proposed and these backgrounds are known to suffer from ghost instabilities [5]. An interesting possibility then is to look for other such solutions in higher codimensional set-ups, initially introduced to address the cosmological constant problem [6,7,9,10,11,12]. This might lead to ghost free models [13] (see however [14]).

In this paper, as a continuation of [15], we consider a codimension two brane that lies at the intersection of two codimension one branes embedded in a six dimensional space. This system was studied in the past in the context of standard gravity [16] (cosmological properties were investigated in [17]), and Gauss-Bonnet gravity [18] elaborating an idea developed in [20]. The latter was generalized to higher-codimensional models in [19]. Models with a generic angle between two intersecting branes were first considered in [21] and then further generalised into the so-called Origami-world in [22]. More recently, in [15], we added brane induced gravity terms to this system to analyse the features of a configuration of
static branes embedded in a time dependent, maximally symmetric background. We showed the existence of new self-accelerating solutions, and of configurations with potentially interesting self-tuning properties.

In the present paper, we continue the analysis of this system by studying cosmological models, obtained by the motion of one of the branes through the bulk, in a mirage approach [23]. The energy momentum tensor different from pure tension on the branes causes the brane to move and bend in the bulk, and induces cosmological evolutions from the point of view of observers sitting on the branes. We allow the branes to intersect at an arbitrary angle and to deform in the preferred shape.

The analysis of gravitational [24] and cosmological [25] aspects of codimension two brane-worlds is a subject that is receiving some attention. Cosmology is mainly studied in the context of a mirage approach. In higher codimensional brane-worlds, the mirage approach has usually some drawbacks (critically examined, for example, in the introduction of [26]), mainly due to fine-tuning relations that the brane energy momentum tensors must satisfy. These are usually associated with the fact that an analogue of Birkhoff theorem does not hold in this case, in contrast to the codimension one case. In codimension one case, this theorem ensures that a system composed by a homogeneous and isotropic brane, moving through a static higher dimensional space, fully catches all the relevant time dependence of the system [27]. In our case, this is not true: generically, a moving higher co-dimensional brane induces time-dependent effects in the bulk [18, 28]. In order to avoid the time dependence in the bulk, one must impose a static ansatz for the bulk geometry, and this is reflected on fine-tuning relations between matter on the brane and in the bulk. Nevertheless, it remains the most direct approach to study cosmological aspects of these models analytically.

The most interesting problem in this system is the isolation of the six dimensional effects in the induced Friedmann equation on the codimension two intersection. As we will see, the Friedmann equation at the intersection receives contributions due to induced gravity terms on it which ensure the recovery of normal 4d cosmology in the relevant regimes. Moreover, there are terms coming from induced gravity on the codimension one branes, of the typical DGP form [2]. Finally, and more interestingly in our framework, the Friedmann equation contains also contributions that come from the six dimensional bulk. They vanish in the limit in which the branes intersect at a right angle, but for generic brane configurations they play an important role for the cosmological evolutions. Indeed, they can provide the late time acceleration, regardless of the energy content of the codimension two brane, generalising the self-accelerating branch of the codimension one DGP model to higher codimensions [2]. This fact has been realized already in [15], but the present analysis is more general because we do not impose the maximal symmetry on the branes under consideration. By properly choosing the embedding for the codimension one branes, the six dimensional effects at the intersection can depend on the inverse of the induced Hubble parameter, and we will analyse the consequences of this in our discussion. Another peculiar feature of our construction is that six dimensional contributions to the Friedmann equation are also associated to the non-conservation of the energy density at the intersection. During the cosmological evolution, the energy density indeed flows from the codimension one to the codimension two branes, unless the codimension one branes intersect with a right angle.

This paper is organised as follow. In Section 2 we will present the general formalism that is necessary to study cosmological properties of the systems we are interested in. In Section 3 we apply this formalism to a particular embedding for the codimension one branes, and in Section 4 we study in some detail cosmological solutions derived from this embedding. Then, in Section 5 we study applications of these cosmological solutions to some interesting situations. We conclude in Section 6.
2 The general formalism

2.1 The model

We consider a system of two intersecting codimension one branes embedded in a six dimensional space-time. They intersect on a four dimensional codimension one brane, where observers like us can be localised. We take an Einstein-Hilbert action for gravity in the bulk and we allow for induced gravity terms on the codimension one branes, as well as on the intersection. Besides gravity, we allow for a cosmological constant term in the bulk, $\Lambda_B$, and for additional fields localised on the branes described by general Lagrangians $L$’s. The general action takes the form

$$S = \int_{\text{bulk}} d^6 x \sqrt{-g} \left( \frac{M_6^4}{2} R - \Lambda_B \right) + \sum_{i=1}^{2} \int_{\Sigma_i} d^5 x \sqrt{-g_{(i)}} \left( \frac{M_{5,i}^3}{2} R_{(i)} + L_{(i)} \right) + \int_{\Sigma \cap} d^4 x \sqrt{-g_{\cap}} \left( \frac{M_4^2}{2} R_{\cap} + L_{\cap} \right),$$

where $\Sigma \cap \equiv \bigcap_i \Sigma_i$ denotes a three-brane at the intersection between all codimension-one branes $\Sigma_i$. We can have different fundamental scales in the different regions of the space, $M_6$, $M_{5,i}$, and $M_4$. The induced gravity terms could be generated, as it was proposed in the original model, by quantum corrections from matter loops on the brane. It is also interesting to note that induced curvature terms appear quite generically in junction conditions of higher codimension branes when considering natural generalisations of Einstein gravity [29, 30] as well as in string theory compactifications [31], orientifold models and intersecting D-brane models [32].

The six dimensional bulk is characterised by a maximally symmetric geometry

$$ds^2 = A^2(t, z^1, z^2) \left( \eta_{\mu
u} dx^\mu dx^\nu + \delta_{kh} dz^k dz^h \right),$$

$$A(t, z^1, z^2) = \frac{1}{1 + \bar{H} t + k_i z^i}.$$  

The parameters $\bar{H}$ and $k_i$ appearing in the warp factor $A$ satisfy the following relation

$$\frac{\Lambda_B}{10} = \bar{H}^2 - k_1^2 - k_2^2,$$  

in order to solve the Einstein equations in the bulk.

We embed a moving and a static codimension one branes ($\Sigma_2$ and $\Sigma_1$ respectively) on the background given by (2). The moving brane $\Sigma_2$ is characterised by an embedding

$$X^{(2)} M = (t, \bar{x}_3, Z_1(t, \omega_1), Z_2(t, \omega_1)).$$

Here, $w_1$ is an embedding coordinate. In the following, for simplicity, we will demand that the intersection with the other brane lies at the position $w_1 = 0$. At this position, we assume that the function $\dot{Z}_1(t, 0)$ vanishes: that is, the second codimension one brane is static at the position of the intersection. The vectors $V$ tangent to $\Sigma_2$ are given by (we introduce indices on the left of $V$: they indicate which brane we are talking about)

$$^{(2)} V^M_{(a)} = \frac{\partial X^M}{\partial x^a}, \quad x^a = (t, x^i, w_1),$$
where \((2) V^{(i)}_M\) is (proportional to) the velocity vector
\[
(2) V^{(i)}_M = (1, 0^i, \dot{Z}_1, \dot{Z}_2) = \dot{X}^M .
\] (6)
The other four vectors are
\[
(2) V^{(i)}_M = (0, \delta^M_i, 0, 0),
\] (7)
\[
(2) V^{(w_{1})}_M = (0, 0^i, Z'_1, Z'_2) = X'^M .
\] (8)
The normal vector to the brane is thus given by the conditions
\[
n_M V^{(a)}_M = 0, \quad \forall a .
\] (9)
Orthogonality with respect to the \(i\) vectors simply removes from \(n_M\) all its 3-dimensional space-like components. Imposing orthogonality w.r.t. \((2) V^{(w_1)}_M\) we then find
\[
n^{(2)}_M = \frac{A}{N} \left( \dot{Z}_1 Z'_2 - \dot{Z}_2 Z'_1, \ddot{0}_3, -Z'_2, Z'_1 \right),
\] (10)
with
\[
N \equiv \sqrt{Z'^2_1 + Z'^2_2 - \left( \dot{Z}_1 Z'_2 - \dot{Z}_2 Z'_1 \right)^2} .
\]
Doing exactly the same steps for the static brane \(\Sigma_1\), with embedding
\[
X^M_{(1)} = (t, \tilde{x}_3, 0, z_2) ,
\] (11)
the vectors tangent to the brane, \((1) V^{(a)}_M\), are immediate to find. And the normal is simply
\[
n^{(1)}_M = A \left( 0, \ddot{0}_3, 1, 0 \right) .
\] (12)

2.2 An useful change of coordinates
Proceeding identically as in the static case\cite{15, 22}, it is useful to change a frame in order to impose the \(Z_2\) symmetry in the case of a general angle. We go to coordinates parallel to the branes \((z_1, z_2) \rightarrow (\tilde{z}_1, \tilde{z}_2)\), where
\[
d\tilde{z}^k \equiv n^{(k)} \cdot dz.
\]
One obtains two two-vectors \(l^{(1)}\) and \(l^{(2)}\) parallel to the branes:
\[
l^{(1)} = \frac{1}{\tilde{Z}_1} (Z'_1, Z'_2), \quad l^{(2)} = \frac{N}{\tilde{Z}_1} (0, 1),
\] (13)
and \(dz = l^{(k)} d\tilde{z}^k\). Then the components of the vectors \((2) V\) parallel to the moving brane become
\[
(2) V^{(w_1)}_M = \frac{\partial \dot{X}^M}{\partial w_1} = \left( 0, \ddot{0}_3, Z'_1, 0 \right),
\] (14)
and
\[(2)\hat{V}_{(0)} = \left(1, \tilde{t}_3, \dot{\bar{z}}_1, \frac{\dot{\bar{Z}}_2 - \dot{Z}_2' \dot{\bar{Z}}_1}{N}\right). \tag{15}\]

Notice that the consistency relation
\[\frac{\partial^2 X_M^{(2)}}{\partial t \partial w_1} = \frac{\partial^2 X_M^{(2)}}{\partial w_1 \partial t},\]
in our case implies the condition
\[\frac{\partial}{\partial w_1} \left(\frac{\dot{\bar{Z}}_1 - \dot{Z}_2' }{N} \right) = 0. \tag{16}\]

The normal to the moving brane becomes in these coordinates
\[\tilde{n}_M^{(2)} = A \left(\frac{\dot{\bar{Z}}_1 - \dot{Z}_2' }{N}, \tilde{t}_3, 0, 1\right). \tag{17}\]

Now, we must take into account that the branes are fixed points of \(Z_2\) symmetries. We focus our analysis on the moving brane \(\Sigma_2\) in order to compute Israel junction conditions at its position: the analysis for the static brane \(\Sigma_1\) can be performed along similar lines. The \(Z_2\) symmetry acting on the static brane \(\Sigma_1\) implies the invariance of the 6d metric under \(\tilde{z}_1 \rightarrow -\tilde{z}_1\), that can be obtained replacing \(\tilde{z}_1 \rightarrow |\tilde{z}_1|\).

After the change of frame, imposing the \(Z_2\) symmetry, the six dimensional metric becomes
\[\tilde{\gamma}_{mn} = \frac{1}{\tilde{Z}_1^2} \begin{pmatrix} \tilde{Z}_1^2 + \tilde{Z}_2'^2 & N \tilde{Z}_2' \text{sgn}(\tilde{z}_1) \\ N \tilde{Z}_2' \text{sgn}(\tilde{z}_1) & N^2 \end{pmatrix}, \tag{18}\]
with the inverse metric
\[\tilde{\gamma}^{mn} = \frac{1}{N^2} \begin{pmatrix} N^2 & -N \tilde{Z}_2' \text{sgn}(\tilde{z}_1) \\ -N \tilde{Z}_2' \text{sgn}(\tilde{z}_1) & \tilde{Z}_1^2 + \tilde{Z}_2'^2 \end{pmatrix}. \tag{19}\]

On the other hand the induced metric is invariant under bulk reparametrisation, and thus reads
\[ds_{5,\Sigma_2}^2 = A^2(t, w_1) \begin{pmatrix} -\left(1 - \dot{\tilde{Z}}_1^2 - \dot{\tilde{Z}}_2^2\right) dt^2 + d\tilde{x}_3^2 + (\tilde{Z}_1'^2 + \tilde{Z}_2'^2) dw_1^2 \\
+ 2(\dot{\tilde{Z}}_1 \tilde{Z}_1' + \dot{\tilde{Z}}_2 \tilde{Z}_2') \text{sgn}(\tilde{z}_1) dt dw_1 \end{pmatrix}. \tag{20}\]

The inverse metric is given by
\[h^{ab} = A^{-2} \begin{pmatrix} \delta^{ij} & 0 \\ 0 & H^{\alpha\beta} \end{pmatrix}, \tag{21}\]
with
\[H^{\alpha\beta} = N^{-2} \begin{pmatrix} - (\tilde{Z}_1^2 + \tilde{Z}_2'^2) & \dot{\tilde{Z}}_1 \tilde{Z}_1' + \dot{\tilde{Z}}_2 \tilde{Z}_2' \\
\dot{\tilde{Z}}_1 \tilde{Z}_1' + \dot{\tilde{Z}}_2 \tilde{Z}_2' & 1 - \dot{\tilde{Z}}_1^2 - \dot{\tilde{Z}}_2^2 \end{pmatrix}, \tag{22}\]
and \(N^2 = \tilde{Z}_1^2 + \tilde{Z}_2'^2 - (\dot{\tilde{Z}}_1 \tilde{Z}_2' - \dot{\tilde{Z}}_2 \tilde{Z}_1')^2\). At the intersection, the four dimensional metric is given by
\[ds_4 = A^2(t)[-(1 - \dot{\tilde{Z}}_1^2 - \dot{\tilde{Z}}_2^2)dt^2 + dx^2]. \tag{23}\]
2.3 Extrinsic curvature

Given all this information, one can compute the components of the extrinsic curvature at the position of the brane \( \Sigma_2 \), using the general formula

\[
K_{mn} = \tilde{V}_{(m)}^M \tilde{V}^N_{(n)} \nabla_M \tilde{n}_N.
\]  

(24)

Since the expression for \( K_{mn} \) is invariant under bulk reparametrisation, to evaluate the right hand side of the previous expression one can use the six dimensional metric in the original frame, or in the frame parallel to the brane.

The calculation of the regular part of \( K_{mn} \) is easier to work out in the original frame. The non vanishing components are the following

\[
K_{00} = \frac{A}{N} \left( \dot{Z}_2 Z_2' - \dot{Z}_1 Z_1' \right) - \frac{A^2}{N} \left( 1 - \dot{Z}_1^2 - \dot{Z}_2^2 \right) \mathcal{K}(w,t),
\]

(25)

\[
K_{w_1 w_1} = \frac{A}{N} \left( Z_2'' Z_2' - Z_1'' Z_1' \right) + \frac{A^2}{N} \left( Z_1' Z_2' + Z_2' Z_1' \right) \mathcal{K}(w,t),
\]

(26)

\[
K_{0 w_1} = \text{sign}(\tilde{z}_1) \frac{A}{N} \left( \dot{Z}_2 Z_1' - \dot{Z}_1 Z_2' \right) + \text{sign}(\tilde{z}_1) \frac{A^2}{N} \mathcal{K}(w,t) \left( \dot{Z}_1 Z_1' + \dot{Z}_2 Z_2' \right),
\]

(27)

\[
K_{ij} = \frac{A^2 \delta_{ij}}{N} \mathcal{K}(w,t).
\]

(28)

where

\[
\mathcal{K}(w,t) = k_1 Z_2' - k_2 Z_1' + H(\dot{Z}_1 Z_2' - \dot{Z}_2 Z_1')
\]

(29)

In addition, the component \( K_{w_1 w_1} \) of the extrinsic curvature may contain terms localised at the intersection due to the presence of the \( sgn \) functions in the six dimensional metric. Let us then look for the singular pieces of the extrinsic curvature

\[
K_{ab}|_{\text{sing}} = \tilde{V}_{(a)}^M \tilde{V}^N_{(b)} \nabla_M n_N|_{\text{sing}}.
\]

(30)

This quantity is much easier to calculate in the tilted reference frame. There are, a priori, two classes of contributions to the singular pieces. The first one comes from partial derivatives acting on \( n_N \), due to the \( sgn \) function included in \( \mathcal{M} \). However such a contribution is proportional to \( n_M N \) itself as the \( sgn \) function only appears in the prefactor of \( n_M N \), and thus vanishes due to the orthogonality between \( n_M N \) and \( V_{(a)}^M \). Then

\[
K_{ab}|_{\text{sing}} = -V_{(a)}^M V_{(b)}^N n_R \Gamma^{R}_{MN}|_{\text{sing}},
\]

(31)

and since \( \Gamma_{MN}^0 \) has no singular part, one is left with \( \Gamma_{MN}^{w_1} \) whose only singular component is \( \Gamma_{\tilde{z}_1 \tilde{z}_1}^{w_1} = g^{22} \partial_i g_{12} = 2 \frac{\tilde{z}_1' \tilde{z}_2^2 + \tilde{z}_2' \tilde{z}_1^2}{N} \delta(\tilde{z}_1) \). Then, in the end, we find that

\[
K_{w_1 w_1}|_{\text{sing}} = -2 A N \frac{Z_2'}{Z_1'} (Z_1'^2 + Z_2'^2) \delta(w_1),
\]

(32)

is the only singular component of the extrinsic curvature.

For the static brane \( \Sigma_1 \), the extrinsic curvature is simply given by

\[
K^a_b = -k_1 \delta^a_b.
\]

(33)
2.4 Junction conditions

The previous expressions for the extrinsic curvature are important in order to obtain the equations that govern the induced cosmology on the brane. They are dictated by the Israel junction conditions

\[ 2 \left[ K_{ab} \right] \equiv 2 \left[ K_{ab} - K h_{ab} \right] = -\frac{1}{M_6} (S_{ab} + S_{ab}^{loc}), \]  

(34)

where \([X] \equiv (X(\Sigma_{2,+}) - X(\Sigma_{2,-}))/2\), while the induced codimension one brane metric is \(h_{ab}\). The extrinsic curvature tensor evaluated on \(\Sigma_2\) is given by \(K_{ab} = h^M_a h^N_b \nabla_M n_N\) with \(K = K^a a\), and energy momentum tensors relative to matter localised on \(\Sigma_2\), appearing on the right hand side of (34), are calculated in the usual way:

\[ S_{ab} = -\frac{2}{\sqrt{-h(2)}} \frac{\delta \left( \sqrt{-h(2)} L(2) \right)}{\delta h_{ab}^{(2)}}, \]  

(35)

\[ S_{ab}^{loc} = -\delta(\Sigma_1)\delta^\mu_\alpha \delta^\nu_\beta \frac{2}{\sqrt{-h(2)}} \frac{\delta \left( \sqrt{-h} L_\cap \right)}{\delta h_{\mu\nu}^\cap} \equiv \delta(\Sigma_1) \sqrt{-h_\cap} \delta^\mu_\alpha \delta^\nu_\beta S_{\mu\nu}. \]  

(36)

In our model the localised energy-momentum tensor also includes contributions from the induced gravity terms. The last quantity \(S_{ab}^{loc}\) denotes energy momentum tensor that is localised on the intersection \(\Sigma_\cap\) between the branes. Notice the presence of the factor \(\sqrt{h_\cap/h(2)}\) that renders the expression covariant with respect to the metric at the intersection.

In the previous discussion, we learned that the only singular term of the extrinsic curvature for the brane \(\Sigma_2\), that is localised at the intersection, is contained in \(K_{w^1 w^1}\). This implies that the six dimensional contributions to the energy momentum tensor must be proportional to the induced metric, \(S_{\mu\nu}^{loc} = f(x^\mu) h_{\cap \mu\nu}\), for some function \(f\). We still do not know whether this function \(f\) is a constant (in which case, it corresponds to pure tension) or not, since we do not know whether the conservation of the energy holds at the intersection or not. The Codazzi equation holds in this case\[1\]

\[ \nabla_a K^a_b = 0 \quad \Rightarrow \quad \nabla_a S^a_b = 0, \]  

(37)

which means that there is no exchange of energy between the bulk and the codimension one branes. But the previous relations may contain singular terms, associated with an exchange of energy between the codimension one and the codimension two branes. This is indeed what generically happens, and we will encounter an example of this phenomenon in Section 3. Singular terms in the first of the previous formulae can appear if \(K^a_a\) has singular pieces, or if some of its components become singular when covariant derivatives act on them. This possibility occurs when the angle between the brane is not right: then the component \(K^a_a\) is non-vanishing at the intersection, and, being proportional to the \(sign\) function (see eq. (27)), it normally generates an additional singular term.

3 Applications

We then consider a system with the static brane \(\Sigma_1\), and the moving brane \(\Sigma_2\). Here, \(\Sigma_2\) is free to move and bend arbitrarily. We assume that the induced gravity term on the moving codimension one brane

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\[ ^1 \text{The RHS of this formula vanishes because both the bulk and the static brane have maximal symmetry}. \]
vanishes: $M_{5,2} = 0$ for simplicity. For this brane, we take an embedding

$$X^M = (t, \vec{x}_3, Z_1(t, w_1), Z_2(t, w_1)), \quad (38)$$

with

$$Z_1 = w_1 \cos \alpha(t, w_1),$$

$$Z_2 = z_2(t, w_1) + w_1 \sin \alpha(t, w_1), \quad (39)$$

where we wrote the two functions $Z_i$ in terms of the auxiliary functions $z_2$ and $\alpha$. We demand that these functions are continuous with respect to the variable $w_1$, at the position of the intersection $w_1 = 0$, and to avoid subtleties related with the reflection symmetry at the intersection we require smoothness conditions $Z_2'(t, 0) = Z_2''(t, 0) = \alpha'(t, 0) = 0$\(^2\). From these definitions, we have

$$\dot{Z}_1 = -\dot{\alpha} w_1 \sin \alpha, \quad Z'_1 = \cos \alpha - w_1 \alpha' \sin \alpha,$$

$$\dot{Z}_2 = \dot{z}_2 + \dot{\alpha} w_1 \cos \alpha, \quad Z'_2 = z'_2 + \sin \alpha + w_1 \alpha' \cos \alpha. \quad (41)$$

Notice that the equation (16) imposes, for this embedding,

$$\dot{\alpha}(t, 0) = 0, \quad (42)$$

that is, the angle does not change at the position of the intersection $w_1 = 0$. Nevertheless, this quantity can change with time away from this point. This is due to our choice of embedding and simplifies the calculations when we focus on the properties of the intersection.

The junction condition on the static co-dimension one brane gives the tension $\lambda_1$ as

$$\lambda_1 = 6M_{5,1}^2(\bar{H}^2 - k_2^2) + 8M_{5,1}^4k_1. \quad (43)$$

The induced metric on the brane $\Sigma_2$ is obtained by plugging the previous expressions in (20). The complete calculation of the cosmological behaviour for the moving codimension one brane is complicated as the brane is inhomogeneous, but it can be obtained straightforwardly from the general formulae (25)–(28). In the next subsection we discuss some of its properties that are useful when comparing them with cosmology on the codimension two brane.

### 3.1 Cosmology on the moving four brane

The cosmological evolution on the moving brane $\Sigma_2$ is complicated by the fact that its induced scale factor and energy momentum tensor must be inhomogeneous, in order to satisfy Israel junction conditions (25)–(28): such conditions require some off-diagonal components of the energy momentum tensor $S^a_b$ to be non-vanishing. In what follows we will only need the explicit form of the codimension-one equations evaluated at the intersection. We thus concentrate on such a limit where the form of the needed energy momentum tensor is the following:

$$S^a_b = \begin{pmatrix} -p_2 & 0 & 0 & 0 & 0 \\ 0 & p_2 & 0 & 0 & 0 \\ 0 & 0 & p_2 & 0 & 0 \\ 0 & 0 & 0 & p_2 & 0 \\ \chi & 0 & 0 & 0 & p_2 \end{pmatrix}, \quad \text{at } w_1 = 0. \quad (44)$$

\(^2\)Asking that only the first derivative vanish at the intersection may be enough to ensure sufficient smoothness to render the system well behaved. In fact later we will briefly mention a situation where a non-vanishing second derivative $z_2''(t, 0) \neq 0$ can turn out to be useful.
Then the junction conditions impose the following relations

\begin{align}
\rho_2 &= -8M_6^4 N^{-1} K, \\
p_2 &= 8M_6^4 (N^{-1} K + \frac{1}{4} N^{-3} A^{-1} z_2 \cos \alpha), \\
\chi &= 2M_6^4 A^{-1} N^{-3} z_2 \sin \alpha \cos \alpha,
\end{align}

where we define \( N = \sqrt{1 - \cos^2 \alpha \frac{z_2}{2}} \) and \( K = k_1 \sin \alpha - (k_2 + \dot{H} z_2) \cos \alpha \).

### 3.2 Cosmology at the intersection

We start from discussing the contributions from the brane \( \Sigma_2 \) to the codimension two brane. At the intersection, characterised by \( w_1 = 0 \), the induced metric is straightforwardly extracted from the five dimensional one and is simply given by

\[ ds_4^2 = A^2(0,t) \left\{ -[1 - \dot{z}_2^2] dt^2 + dx_3^2 \right\} = -d\tau^2 + a(\tau)^2 d\bar{x}_3^2, \quad d\tau = A(0,t) \sqrt{1 - \dot{z}_2^2} dt. \]

Then the induced Hubble parameter is

\[ H = \frac{1}{a(\tau)} \frac{da(\tau)}{d\tau} = \frac{\ddot{H} + k_2 \dot{z}_2}{\sqrt{1 - \dot{z}_2^2}}. \]

In order to find the Friedmann equation at the intersection, we have to extract the singular part of the Israel junction conditions for the codimension two branes. This singular part receives contributions from the energy momentum tensor localised on the codimension two bran (containing also the induced gravity terms at the intersection), from the induced gravity terms on the codimension one branes, and from singular contributions of the extrinsic curvature terms. The final contribution represents the most interesting feature of our model since it corresponds to six dimensional contributions to four dimensional physics. We start our discussion with their evaluation.

The only singular part on the extrinsic curvature for brane \( \Sigma_2 \) is contained in the \((w_1, w_1)\) component. It is

\[ (K_{w_1}^w)^{\text{sing}} = -\frac{2}{A \sqrt{1 - \cos^2 \alpha \frac{z_2}{2}}} \frac{1 - \dot{z}_2^2}{1 - \cos^2 \alpha \frac{z_2}{2}} \frac{\sin \alpha}{\cos \alpha} \delta(w_1), \]

so that \((\dot{K}_{mn})^{\text{sing}} = -g^{(5)}_{mn} (K_{w_1}^w)^{\text{sing}}\). Then one obtains the following contribution to the energy momentum at the four-dimensional intersection:

\[ T_{\mu\nu}^{\text{loc}} = 2M_6^4 \delta_\mu^n \delta_\nu \frac{\sqrt{h(2)}}{\sqrt{h(\cap)}} (\dot{K}_{mn})^{\text{sing}} = 4M_6^4 \frac{\sqrt{1 - \dot{z}_2^2}}{1 - \cos^2 \alpha \frac{z_2}{2}} \frac{\sin \alpha}{\cos \alpha} h_{\cap \mu\nu}. \]

As expected, it is proportional to the induced metric at the intersection.

Proceeding with our calculation, it is easy to find the contributions to the intersection from the induced gravity terms on \( \Sigma_1 \) (recall that we have chosen \( M_{5,2} = 0 \) so there are no induced gravity terms on the moving brane \( \Sigma_2 \)). We find

\[ (G_0^0)^{\text{sing}} = -\frac{6k_2}{\sqrt{1 - \dot{z}_2^2}}. \]

Putting all this information together, we find the following Friedmann equation

\[ \rho = 3M_4^2 H^2 + 6M_5^2 \frac{k_2}{\sqrt{1 - \dot{z}_2^2}} + 4M_6^4 \frac{\sin \alpha}{\cos \alpha} \frac{\sqrt{1 - \dot{z}_2^2}}{1 - \cos^2 \alpha \frac{z_2}{2}}. \]
The induced gravity terms on the codimension one branes do not induce a violation of the continuity equation at the intersection because the intersection can be seen as codimension one object from the point of view of the four branes and then the properties of the Israel formalism for the codimension one brane ensure the conservation of energy (see eq. (37)). On the other hand, the last, six dimensional term in Eq. (53) is explicitly time dependent, while we know that it appears as a tension term in the effective energy momentum tensor at the intersection. This is because it is proportional to the induced metric (see Eq. (51)). This indicates that this term is likely to be associated with a violation of the continuity equation at the intersection.

This issue can be understood by re-considering the Codazzi equation:

$$\nabla_M \hat{K}^M_N = \left( \nabla_M \hat{K}^M_N \right)^{(\text{reg})} + \left( \nabla_M \hat{K}^M_N \right)^{(\text{sing})} = 0,$$

where $\nabla$ is the covariant derivative with respect to the five dimensional metric on the codimension one brane. From the previous formula, we learn that both the regular and singular parts must vanish simultaneously. However, it can happen that the covariant derivative induces singular contributions when it is applied to certain components of $\hat{K}^M_N$ by taking derivatives of sign functions. This is indeed what happens in our case. Consider the case $N = 0$. The singular part of the previous formula tells us that

$$\frac{\partial}{\partial t} \left( K^{w_1}_{w_1} \right)^{\text{sing}} = \left( \nabla_M K^M_0 \right)^{\text{sing}},$$

where the left hand side contains the singular term of eq. (50), while in the right hand side the singular terms come from the covariant derivatives. But the piece in the left hand side corresponds precisely to the term associated with the six dimensional contribution at the intersection. Thus the six dimensional contribution to the Friedmann equation on the intersection does not satisfy the energy conservation and there is an exchange of energy from codimension two brane to the higher dimensional space.

In the light of this fact, one expects that the conservation of energy at the intersection does not hold. Instead, one finds the continuity equation

$$\dot{\rho} + 3H (\rho + p) = 4M_6^4 \tan \alpha \frac{\partial}{\partial \tau} \left[ \frac{\sqrt{1 - \dot{z}_2^2}}{1 - \cos^2 \alpha \dot{z}_2^2} \right].$$

The conservation of energy is ensured only when $\alpha$ vanishes, or when $\dot{z}_2$ is constant.

We close this section by summarising the equations that govern cosmology of the codimension one branes and the intersection;

$$\frac{\Lambda_B}{10} = \bar{H}^2 - k_1^2 - k_2^2,$$

$$\lambda_1 = 6M_5^2 (\bar{H}^2 - k_2^2) + 8M_6^4 k_1,$$

$$\rho_2 = -8M_6^4 \frac{1}{\sqrt{1 - \cos^2 \alpha \dot{z}_2^2}} \left( k_1 \sin \alpha - (k_2 + \bar{H} \dot{z}_2) \cos \alpha \right),$$

$$\chi = 2M_6^4 \frac{\dot{z}_2 \ddot{z}_2 \sin \alpha \cos \alpha}{A(0, t)(1 - \cos \alpha \dot{z}_2^2)^{3/2}},$$

$$\rho = 3M_6^2 \bar{H}^2 + 6M_{5,1}^2 \frac{k_2}{\sqrt{1 - \dot{z}_2^2}} + 4M_6^4 \tan \alpha \frac{\sqrt{1 - \dot{z}_2^2}}{1 - \cos^2 \alpha \dot{z}_2^2},$$

where $A(0, t) = 1/(1 + \bar{H} t + k_2 \dot{z}_2(t))$, $\Lambda_B$ is the bulk cosmological constant, $\rho_2$ is the energy density on the moving codimension one brane at $w_1 = 0$, $\chi$ is $(w_1, t)$-component of energy momentum tensor on the
moving codimension one brane at \( w_1 = 0 \) and \( \rho \) is the energy density at the intersection. The Hubble parameter at the intersection is given by Eq. (49). The energy conservation at the intersection is given by Eq. (56). In the following we put \( M_{1,5} = M_5 \).

4 Cosmological solutions

In this section, we discuss the property of the cosmological solutions by focusing on the Friedmann equations on the moving four brane and at the intersection.

4.1 The branes at a right angle

We first consider the simplest case in which the branes are at a right angle \( \alpha = 0 \) and \( \bar{H} = 0 \). In this case, the energy momentum tensor on the moving four brane becomes the perfect fluid and there is no energy flow \( \chi = 0 \). Using the cosmic time \( \tau \), the 5D metric is given by

\[
 ds_5^2 = -d\tau^2 + A^2(w_1, t) \left( d\vec{x}_3^2 + dw_1^2 \right), \quad A(w_1, \tau) = \frac{1}{1 + k_1 w_1 + k_2 z_2(\tau)}. \tag{62}
\]

Although the scale factor depends on \( w_1 \), the Hubble parameter in terms of the cosmic time is independent of \( w_1 \) and given by Eq. (49). By expressing \( \dot{z}_2 \) in terms of the Hubble parameter \( H \), we get

\[
 \rho_2 = 8M_6^4 \sqrt{H^2 + k_2^2}. \tag{63}
\]

Since there is no energy flow, the energy density is conserved

\[
 \partial_\tau \rho_2 + 4H(\rho_2 + p_2) = 0. \tag{64}
\]

At the intersection, the Friedmann equation is given by

\[
 \rho = 3M_2^2 H^2 + 6M_5^3 k_2 \sqrt{1 + \frac{H^2}{k_2^2}}, \tag{65}
\]

and the standard continuity equation holds since \( \tan \alpha = 0 \). Notice that the static brane gives a contribution of the 5D DGP form.

Since the Hubble parameters are equal in both the equations (63) and (65), by expressing \( H \) as a function of the energies in the two cases and equalling the results, they will imply a fine tuning relation between the two homogeneous energy densities \( \rho_2 \) and \( \rho \). Geometrically, this is because when the codimension one brane \( \Sigma_2 \) moves through the static bulk, it completely controls the dynamics of the brane \( \Sigma_\cap \) that sits at the intersection with \( \Sigma_1 \). Then, \( \Sigma_\cap \) can only follow the motion of \( \Sigma_2 \), without an independent dynamics on its own. The problem becomes clearer by the fact that the energy density and the Hubble parameter on the moving codimension brane do not depend on the coordinate \( w_1 \). Then, the energy density at the intersection actually fixes all the properties of the energy density on the moving brane \( \Sigma_2 \), including its equation of state.

There is a simple way out of a part of this problem. The fine-tuning we found is so strong because we demand that the moving brane \( \Sigma_2 \) keeps a straight shape – that is, it cannot deform along the \( z_1 \) direction. Suppose however that we allow the moving codimension one brane to be free to deform and bend, forming a non-trivial angle \( \alpha \) with \( \Sigma_1 \) which explicitly depends both on \( z_1 \) and \( t \). Then, the energy density on \( \Sigma_2 \) will explicitly depend on \( z_1 \). This implies that, although the energy density at the intersection must equal the energy density on \( \Sigma_2 \) calculated at \( z_1 = 0 \), nevertheless this fine-tuning is ameliorated with respect to the previous case. Indeed, it involves only the quantities calculated at the intersection.
4.2 Arbitrary angle between the branes

Now let us consider the case in which $\alpha \neq 0$. In this case, the Hubble parameter at $w_1 = 0$ is given by Eq. (49). However, it is important to emphasise that this is only the value of the Hubble parameter when evaluated at the position of the intersection: when calculated away from this point, it receives additional contributions coming from time derivatives of the angle $\alpha$, and it becomes an inhomogeneous quantity. This is a crucial difference with respect to the example studied in the previous subsection. Again taking $\bar{H} = 0$ for simplicity, the Friedmann equation on the moving four brane at the position of the intersection is given by

$$\rho_2 = 8M_6^4 (k_2 \cos \alpha - k_1 \sin \alpha) \sqrt{\frac{H^2 + k_2^2}{H^2 \sin^2 \alpha + k_2^2}}. \quad (66)$$

On the other hand, the Friedmann equation at the intersection is given by

$$\rho = 3M_4^2 H^2 + 6M_6^2 k_2 \sqrt{1 + \frac{H^2}{k_2^2}} + 4M_6^5 k_2 \tan \alpha \left[ \sqrt{\frac{H^2 + k_2^2}{H^2 \sin^2 \alpha + k_2^2}} \right]. \quad (67)$$

When the angle $\alpha$ vanishes and $k_2$ remains finite, we recover the results of the previous subsection. Notice also that comparison between Eqs. (66) and (67) imposes a fine-tuning relation between $\rho$ and the energy density $\rho_2$ of the codimension one brane, when evaluated at the intersection. Nevertheless, this fine-tuning is much milder than the one we met in the previous subsection. This fine-tuning is associated with the restrictive Ansatz we have chosen for the bulk metric.

The form of the previous Friedmann equation is quite complicated to study with full generality. While the second term on the right hand side of Eq. (67) corresponds to the well-known DGP-like term, the last term in the right hand side of Eq. (67) is less standard, and is associated with six dimensional contributions. This term is interesting because they contain $H$ at the denominator in a peculiar way, with interesting consequences for cosmology. The fact that $H$ appears at the denominator implies that for large $H$ this term is suppressed. One may be tempted to interpret this behaviour, at least partially, as a relativistic effect. Indeed, large $H$ means that the brane speed is approaching the speed of light (see the definition in formula (49)). But at higher and higher speeds, due to the Lorentz contraction, an observer on the moving brane sees the intersection angle with the static brane approaching the value $\pi/2$. But we know that in this limit the six dimensional contributions at the intersection vanish, explaining why for large $H$ this term is suppressed.

The continuity equation is given by

$$\dot{\rho} + 3H (\rho + p) = 4M_6^5 k_2 \tan \alpha \frac{\partial}{\partial t} \left[ \frac{\sqrt{H^2 + k_2^2}}{H^2 \sin^2 \alpha + k_2^2} \right], \quad (68)$$

so, under the assumption that $k_2$ is non zero, the conservation of energy is ensured only when $\alpha$ vanishes, or when $H$ is constant. For $\alpha \neq 0$, it is also necessary to have the energy flow on the moving codimension one brane, $\chi$, which also breaks the conservation of energy on the moving codimension one brane. Then we can understand that the energy at the intersection is transmitted to the moving codimension one brane.

5 Applications

In this section, we derive some interesting consequences of the cosmological equations we discussed in the previous section. In the first two subsections, we examine some of the cosmological properties of the
solutions we discussed in [15] in the present context. In the last subsection, we will instead derive a new selfaccelerating configuration in which the codimension one branes are not maximally symmetric, a case that we did not discuss in our previous work.

5.1 Self-tuning solutions

Here we re-examine the self-tuning solution presented in [15]. We take \( k_1 = 0, k_2 = 0, M_{5,1} = 0 \) and \( \dot{z} = 0 \). Then we have

\[
\lambda_4 = 3M_4^2 \bar{H}^2 + 4M_6^4 \tan \alpha, \quad \bar{H}^2 = \frac{\Lambda_B}{10}. \tag{69}
\]

where \( H = \bar{H} \). Then the expansion rate does not depend on \( \lambda_4 \). The self-tuning mechanism consists on the fact that if we change tension \( \lambda_4 \), \( \alpha \) changes so that induced cosmology remains the same. Unfortunately, our embedding is not well suited to study the self-tuning property of the solution as \( \alpha = \text{constant} \) is imposed from the consistency relation (16), as we find in equation (42). A possible way out would be to consider a situation in which \( z''(t,0) \neq 0 \) at the intersection. This would generalise (16) with new pieces that would not necessarily impose that \( \dot{\alpha}(t,0) = 0 \). It would be nice to study in more detail this kind of generalisation to understand whether it can be compatible with the reflection symmetries of our system or not. It is important to understand whether the eventual self-tuning property would be compatible with the recovery of small scale 4d general relativity on the intersection, in order not to contradict big bang nucleosynthesis and other cosmological tests. In order for this last tricky issue to be solved, it seems to be necessary that the dynamical angle reacts only to the vacuum energy density component of the localised matter on the intersection: a priori this is rather counterintuitive. However, a few observations are in order here: it is well known that 6d brane worlds with conical singularities treat tension-type of matter on a completely different footing with respect to a generic fluid (\( \omega \neq -1 \)) [34]. In fact, also in the present setup the six-dimensional contribution to the 4d stress tensor (cfr. eq. (51)) has a tension-like structure; moreover as we showed, a generic fluid localised on the intersection does not seem to render the bulk geometry singular as opposed to what happens in the thin conical setups [34], but violates the conservation of energy. It is therefore not excluded that the self-tuning might be at work here.

5.2 Self-accelerating solutions with \( \bar{H} \neq 0 \)

Here we re-consider the self-accelerating solution presented in [15], for maximally symmetric configurations. We assume there is no cosmological constant nor matter in the system \( \Lambda_B = \rho = \rho_2 = \lambda_1 = 0 \) with \( \dot{z} = 0 \). Then we get

\[
0 = \bar{H}^2 - k_1^2 - k_2^2, \tag{70}
\]
\[
0 = 3M_4^2 \bar{H}^2 + 6M_5^3 k_2 + 4M_6^4 \tan \alpha, \tag{71}
\]
\[
0 = k_1 \sin \alpha - k_2 \cos \alpha, \tag{72}
\]
\[
0 = 6M_5^3 (\bar{H}^2 - k_2^2) + 8M_6^4 k_1. \tag{73}
\]

If \( k_1 < 0 \) and \( k_2 < 0 \), there are non-trivial solutions for \( k_1, k_2, \alpha \) and \( \bar{H} \). The solution is roughly given by

\[
\bar{H} \sim \frac{M_2^2}{M_4}, \quad M_5^3 \sim M_4 M_6^2. \tag{74}
\]

in accordance with what we found in the previous paper. Once \( \alpha, k_2 \) and \( \bar{H} \) are fixed, Eqs. (56) and (61) determine the cosmological dynamics with \( \rho \) without ambiguity. The resulting cosmology is complicated due to the non-conservation of energy. Instead of dealing with this complicated case, we will discuss a simpler, different situation in the next subsection.
5.3 Self-accelerating solution with $\bar{H} = 0$

We consider the case $\bar{H} = 0$, $k_1 = 0$, and $k_2^2 = \beta^2 \sin^2 \alpha M_6^2$ for some constant $\beta$ that we take very small. Then, we regard the quantity $H^2$ as much bigger than $\beta^2 M_6^2$ (we will see that this approximation can be satisfied in our context). The continuity equation becomes

$$\dot{\rho} + 3H (\rho + p) = -4\epsilon \frac{\beta M_6^5}{\cos \alpha} \frac{\partial}{\partial t} \frac{1}{H},$$

(75)

while the Friedmann equation acquires the form

$$\rho \simeq 3M_4^2 H^2 - 6M_5^3 \epsilon H - 4\epsilon \frac{\beta M_6^5}{\cos \alpha H},$$

(76)

with $\epsilon \equiv -\frac{\sin \alpha}{\sin \alpha} = \pm 1$. The first term in the right hand side is dominant at large $H$, and ensures the correct four dimensional form for early time cosmology. The second term is the typical DGP contribution, while the third term, a six dimensional effect, is less standard as we discussed before. The previous expression can be easily rewritten as

$$\frac{\rho}{3M_4^2} = \left( H - \epsilon \frac{M_5^3}{M_4^2} \right)^2 - \frac{M_6^5}{M_4^4} - 4\epsilon \frac{\beta M_6^5}{3M_4^2 \cos \alpha H},$$

(77)

from which we obtain

$$H = \epsilon \frac{M_5^3}{M_4^2} + \sqrt{\frac{\rho}{3M_4^2} + \frac{M_6^5}{M_4^4} + 4\epsilon \frac{\beta M_6^5}{3 \cos \alpha M_4^2 H}},$$

(78)

by imposing that the quantity inside the square root is positive. Now, the choice $\epsilon = 1$ corresponds to the standard DGP self-accelerating branch, and the quantity inside the square root is always positive. It implies that, even when $\rho$ vanishes, the Hubble parameter satisfies the inequality

$$H \geq \frac{M_5^3}{M_4^2} \left( 1 + \sqrt{1 + 4 \frac{\beta M_6^5 M_4^2}{M_5^3 \cos \alpha H}} \right),$$

(79)

and so we find a lower bound for $H$, as in the well-known self-accelerating branch of DGP model.

This case is very similar to the standard five dimensional case, since the acceleration is mainly driven by the effects of the codimension one brane. It is however also possible to study the case in which $M_5 = 0$, to understand whether six dimensional effects provide acceleration by themselves. Then, the continuity equation (75) can be formally integrated as

$$\rho = \rho_0 \left( \frac{a(t)}{a_0} \right)^{-3(1+w)} \left( 1 - \frac{4\beta M_6^5}{3H^3 M_4^2 \cos \alpha} \right)^{\frac{3}{2}},$$

(80)

with two constants $\rho_0$ and $a_0$ where $w = p/\rho$. This solution shows that, in the limit $H \to \infty$, one recovers the usual relation between energy density and scale factor. Plugging this expression in the Friedmann equation, a little manipulation leads to the following relation

$$\frac{3}{2} \rho_0 \left( \frac{a(t)}{a_0} \right)^{-\frac{3}{2}(1+w)} = (3M_4^2)^{\frac{3}{2}} \left( H^3 - \frac{4\beta M_6^5}{3M_4^2 \cos \alpha} \right),$$

(81)
that is the differential relation that our scale factor must satisfy, for a given equation of state. We notice that from this one gets the following relation for the acceleration:

\[
\frac{\ddot{a}}{a} = \dot{H} + H^2 = \left[ 2 - 3 \left( 1 + \omega \right) \left( 1 - \frac{4\beta M_6^5}{3 H^3 M_4^2 \cos \alpha} \right) \right] \frac{H^2}{2},
\]

so we have the acceleration when

\[
\omega < \frac{2}{3 \left( 1 - \frac{4\beta M_6^5}{3 H^3 M_4^2 \cos \alpha} \right)} - 1.
\]

Then for small \( H \) we learn that the six dimensional contributions help to provide the acceleration and it can be achieved even when \( \omega > -\frac{1}{3} \). At late times \( \rho \to 0, H \to (4\beta M_6^5 / 3M_4^2 \cos \alpha)^{1/3} \) and the solution approaches de Sitter solution. This cosmological model can be studied along the lines of the analysis of [35]. Notice that the present example of self-acceleration is different with respect to the ones discussed in the previous subsection originally found in [15]. This is because, for the particular choice of our embedding, the codimension one branes do not need to be maximally symmetric, nor empty. It would be nice to understand whether in this case ghosts are present in the low energy spectrum, and if so how do they manifest themselves.

### 6 Conclusions and Open Issues

In this paper, we explored the cosmological features of a codimension two brane with induced gravity terms, sitting at the intersection between two codimension one branes in six dimensions. We found that the cosmological expansion at the intersection is controlled by contributions coming from the codimension one branes, and from the six dimensional bulk. We first showed that the effect of the codimension one branes on the Friedmann equation at the intersection have the well-known DGP form. Then, we learned that six dimensional contributions are much less standard. They can have an important role for late time cosmology providing a new source of the geometrical acceleration, controlled by the angle between the branes. At the same time, they are also associated with a violation of the energy conservation at the intersection, allowing a flow of energy between the codimension two brane and the higher dimensional space. We discussed consistency relations that matter on the codimension two brane must satisfy and the connection with the choice of energy momentum tensor localised on the codimension one branes.

The main aim of this work was to formulate a general and powerful formalism based on the approach of mirage cosmology that can be used to study cosmological solutions in this and similar models, and to apply it to a couple of representative examples. Due to the fact that the codimension one branes that intersect with general angle are not homogeneous and isotropic, a numerical analysis is likely to be needed in order to analyse in full details the cosmological evolution of this class of models. As a natural continuation of the present work, it would be interesting to study the low energy effective action for the light modes associated with our brane configurations. This analysis would be necessary in order to investigate whether ghosts are present in the spectrum of the low energy theory, and, if so, whether they can be eliminated with a mechanism similar to the one of [13]. We leave these issues to future work.

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