Cosmological Acceleration from Energy Influx

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Abstract

We discuss the cosmological evolution of a 3-brane Universe in the presence of energy influx from the bulk. We show that this influx can lead to accelerated expansion on the brane, depending on the equations of state of the bulk and brane matter. The absorption of non-relativistic bulk matter by the brane at an increasing rate leads to a small positive acceleration parameter during the era of matter domination on the brane. On the other hand, the brane expansion remains decelerating during radiation domination.
**Introduction:** Recent astronomical observations indicate that the expansion of our Universe is accelerating [1]. The physical mechanism that drives this cosmological acceleration has not been established yet. The most popular explanation relies on the presence of a small cosmological constant or dark energy. However, in order not to disturb successful ingredients of standard Big Band cosmology such as nucleosynthesis, the contribution of the dark energy to the total energy density must be strongly constrained. As a result, the dark energy can be significant only during recent times. A convincing explanation of this “cosmic coincidence” has not been given yet. Most scenaria identify the dark energy with the vacuum energy of a time-dependent field with carefully chosen dynamics [2].

In this letter we would like to explore another possibility, that does not require the presence of a field. The mechanism can be demonstrated by considering the equation

$$\dot{\rho} + 3H(\rho + p) = T,$$

(1)

that describes energy conservation in an expanding Universe with Hubble parameter $H$, in the presence of a source term $T$. The standard conservation equation with $T = 0$ has stationary solutions with $\dot{\rho} = 0$ only if $\rho + p = 0$. This implies the presence of vacuum energy with equation of state $p = -\rho$, that leads to the standard inflationary scenario with constant $H$ and an exponentially increasing scale factor $a$.

However, if $T \neq 0$ stationary solutions are possible for a more general class of equations of state. For example, one could have constant $H, \rho, p$, such that $3H(\rho + p) = T$. The resulting cosmological solution has an exponentially increasing scale factor and positive acceleration. It has been termed “steady-state Universe” [3]. Interesting behavior, with positive acceleration, can be obtained even if $T$ is not constant, but has some dependence on the energy density $\rho$ or the scale factor $a$.

In the framework of a four-dimensional world, the energy source term $T$ does not have a simple origin. However, if more than four dimensions exist, our observable Universe may be identified with a 3-brane embedded in the higher-dimensional world. In this case, eq. (1) would refer to the matter localized on the brane. The source term $T$ could have the simple interpretation of energy density falling onto the brane from the extra dimensions. The energy-momentum tensor in the higher-dimensional world is conserved and no source terms exist. However, an effective source term is generated for the matter localized on the brane. In the following we shall discuss how such a scenario can be realized.

**The framework:** We work in the framework of the Randall-Sundrum (RS) model [4]. It is described by the action

$$S = \int d^5x \sqrt{-g} \left( M^3R - \Lambda + L_{B}^{\text{mat}} \right) + \int d^4x \sqrt{-\hat{g}} \left( -V + L_b^{\text{mat}} \right),$$

(2)

where $R$ is the curvature scalar of the five-dimensional metric $g_{AB}, A, B = 0, 1, 2, 3, 4$, and $\hat{g}_{\alpha\beta}$, with $\alpha, \beta = 0, 1, 2, 3$, is the induced metric on the 3-brane. $\Lambda$ is the bulk cosmological constant, while the quantity $V$ includes the brane tension as well as quantum contributions to the four-dimensional cosmological constant. We consider the commonly used ansatz for the metric [5, 6]

$$ds^2 = -n^2(t, z)dt^2 + a^2(t, z)\gamma_{ij}dx^idx^j + b^2(t, z)dz^2,$$

(3)

where $\gamma_{ij}$ is a maximally symmetric 3-dimensional metric. We use $k = -1, 0, 1$ to parametrize the spatial curvature.
We decompose the energy-momentum tensor into vacuum and matter contributions in the bulk and on the brane

\[
T^A_C = T^A_C|_{\nu,b} + T^A_C|_{\nu,b} + T^A_C|_{\nu,B} + T^A_C|_{\nu,B}
\]  

(4)

\[
T^A_C|_{\nu,b} = \frac{\delta(z)}{b}\text{diag}(-V,-V,-V,0)
\]  

(5)

\[
T^A_C|_{\nu,B} = \text{diag}(-\Lambda,-\Lambda,-\Lambda,-\Lambda,-\Lambda)
\]  

(6)

\[
T^A_C|_{m,b} = \frac{\delta(z)}{b}\text{diag}(-\rho,\rho,\rho,\rho,0),
\]  

(7)

where \(\rho\) and \(p\) are the energy density and pressure on the brane, respectively. The behaviour of \(T^A_C|_{m,B}\) is in general complicated in the presence of flows, but we do not have to specify it further in this work.

We are interested in the Einstein equations at the location of the brane. We indicate by the subscript \(o\) the value of various quantities on the brane. It is convenient to work in a coordinate frame in which \(b_o = n_o = 1\). This can be achieved by using coordinates such that \(b(t, z) = 1\). This is always possible as can be checked by considering the “two-dimensional” part of the metric: \(-n^2(t, z) dt^2 + b^2(t, z) dz^2\). Through an appropriate coordinate transformation we can always set this part in the form \(-n^2(t, z) dt^2 + dz^2\). We can then redefine the time coordinate so that \(n_o = 1\) on the brane. We emphasize that our assumptions for the form of the energy-momentum tensor become now specific to this coordinate frame. In particular, the brane is identified with the hypersurface \(z = 0\) in this frame. In this sense, our discussion becomes more restricted. The gained advantage is that significant progress can be made without solving explicitly the Einstein equations in the bulk in the presence of flows.

**The Einstein equations:** For \(b(t, z) = n_o = 1\) we find \([5]–[8]\)

\[
\dot{\rho} + 3\frac{\dot{a_o}}{a_o}(\rho + p) = -2T^0_4
\]  

(8)

\[
\frac{\ddot{a_o}}{a_o} + \left(\frac{\dot{a_o}}{a_o}\right)^2 + \frac{k}{a_o^2} = \frac{1}{6M^3}\left(\Lambda + \frac{1}{12M^3}V^2\right) - \frac{1}{144M^6}(V(3p - \rho) + \rho(3p + \rho)) - \frac{1}{6M^3}T^4_4,
\]  

(9)

where \(T^0_4, T^4_4\) are the 04 and 44 components of \(T^A_C|_{m,B}\) evaluated on the brane.

We are interested in a model that reduces to the RS vacuum [4] in the absence of matter. In this case, the first term in the r.h.s. of eq. (9) vanishes. A new scale \(k_{RS}\) is defined through the relations \(V = -\Lambda/k_{RS} = 12M^3k_{RS}\).

We can rewrite eqs. (8), (9) in the equivalent form \([7, 8]\)

\[
\dot{\rho} + 3H_o(\rho + p) = -2T^0_4
\]  

(10)

\[
H_o^2 = \frac{\dot{a_o}^2}{a_o^2} = \frac{1}{144M^6}\left(\rho^2 + 2V \rho\right) - \frac{k}{a_o^2} + \chi + \phi + \lambda
\]  

(11)

\[
\dot{\chi} + 4H_o \chi = \frac{1}{36M^6}(\rho + V)T^0_4,
\]  

(12)

\[
\dot{\phi} + 4H_o \phi = \frac{-1}{3M^3}H_o T^4_4.
\]  

(13)

The conformally flat form \(e^2(z, t)(-dt^2 + dz^2)\) is more commonly used in the literature. The important point is that, using the reparametrization invariances, a two-dimensional metric can be put in a form that depends only on one function of \(t, z\).
where $\lambda = (\Lambda + V^2/12M^3)/12M^3$ is the effective cosmological constant on the brane. The functions $\chi, \phi$ are defined through eqs. (12), (13). In the RS model $\lambda$ is zero, the value we shall use in the rest of the paper. We have assumed the orbifold symmetry $z \leftrightarrow -z$, so that $2T^0_4$ is the discontinuity of the 04 component of the bulk energy-momentum tensor at the location of the brane.

In the low-density region, in which $\rho \ll V$, we may ignore the term $\sim \rho$ in the above equations compared to $V$ and define $M^2_{Pl} = 12M^6/V = M^3/k_{RS}$. The ratio of the terms in the r.h.s. of eqs. (13), (12) is of order $T^4_0 T^4_0 M^3 H^2_o \sqrt{\rho} V \sim T^4_0 T^4_0 \sqrt{\rho} V$. (14)

We assume that $\rho/V$ is sufficiently small for this ratio to be small at all times of interest\(^2\). Then $\phi$ can be set to zero and omitted from our considerations. As a result, eqs. (10)–(12) can be written as

$$\dot{\rho} + 3(1 + w) H_o \rho = -2T^0_4$$

(15)

$$H^2_o = \left(\frac{\dot{a}_o}{a_o}\right)^2 = \frac{\rho}{6M^2_{Pl}} + \chi - \frac{k}{a^2_o}$$

(16)

$$\dot{\chi} + 4H_o \chi = \frac{1}{3M^2_{Pl}} T^0_4,$$

(17)

where $p = w \rho$. The cosmological evolution is determined by three initial parameters ($\rho_i, a_i, \chi_i$, or alternatively $\rho_i, a_i, \dot{a}_i$), instead of the two ($\rho_i, a_i$) in conventional cosmology. The reason is that the generalized Friedmann eq. (11) (or (16)) is not a first integral of the Einstein equations because of the possible energy exchange between the brane and the bulk. In the above equations, we recover the “mirage” or “Weyl radiation” component $\chi$, found in studies of the cosmological evolution in the presence of energy outflow [8, 9].

The bulk energy-momentum tensor: The determination of the cosmological evolution on the brane requires information on the form of $T^0_4$. This can be determined exactly only through the solution of the Einstein equations in the bulk, a formidable task in the case of energy flows. However, the qualitative behaviour of $T^0_4$ can be inferred from the form of the energy-momentum tensor. For a perfect fluid with five-velocity $U^A$ it has the form

$$T^{AC} \big|_{m,B} = p_B g^{AC} + (p_B + \rho_B) U^A U^C.$$  

(18)

For a fluid falling onto the brane with velocity $v_4$ along the fifth dimension, and within our ansatz for the metric, we have $U^0 = (1 - v^2_4)^{-1/2}$, $U^i = 0$, $U^4 = v_4 U^0$ at the location of the brane. As there is no expansion along the fifth dimension, we expect that the main effect of the expansion along the remaining three spatial dimensions is to dilute the fluid energy density and reduce the pressure. These should fall with a certain power of the scale factor, determined by the equation of state of the fluid.

\(^2\)It is obvious that the omission of $\phi$ is not justified when the energy influx stops and $T^0_4$ becomes zero. However, we assume that at this point $T^0_4$ is also very small, so that $\phi$ gives a negligible contribution to eq. (11).
In order to determine this power we may consider the conservation of the energy-momentum tensor in the absence of energy flows. For the 00 component of $T^A_{C|m,B}$ at the location of the brane we find
\[ \dot{T}^0_0 = \frac{\dot{a}_o}{a_o} \left( -3 T^0_0 + T^i_i \right), \] (19)
where we have made use of $b_o = n_o = 1$. For a non-relativistic gas of bulk particles with $T^A_{C|m,B} = \text{diag}(\rho_B, 0, 0, 0, 0)$, we obtain $\rho_B \sim a_o^{-3}$. For a relativistic gas with $T^A_{C|m,B} = \text{diag}(\rho_B, \rho_B/4, \rho_B/4, \rho_B/4, \rho_B/4)$, we obtain $\rho_B = 4p_B \sim a_o^{-15/4}$. For a gas with zero pressure along the fifth dimension and $T^A_{C|m,B} = \text{diag}(\rho_B, \rho_B/3, \rho_B/3, \rho_B/3, 0)$, we obtain $\rho_B = 3p_B \sim a_o^{-4}$. Finally, for the case of vacuum energy with $T^A_{C|m,B} = \text{diag}(\rho_B, -\rho_B, -\rho_B, -\rho_B, -\rho_B)$, we obtain $\rho_B = -p_B \sim \text{const}$.

In analogy with the above, for a non-relativistic gas of bulk particles with constant $v_4$ we expect $T^0_4 \sim a_o^{-3}$ at the location of the brane, while for an isotropic relativistic gas $T^0_4 \sim a_o^{-15/4}$. If $T^A_{C}$ originates in the vacuum energy of a bulk scalar field we expect $T^0_4 \sim \text{const}$. Other types of behaviour are also possible in scenarios in which the velocity $v_4$ varies with time.

We parametrize the dependence of $T^0_4$ on $a_o$ as
\[ \frac{1}{3M_{Pl}^2} T^0_4(t) = \frac{1}{3M_{Pl}^2} T^0_4(t) \left( \frac{a_o(t)}{a_o(t)} \right)^q = -\frac{T}{(a_o(t))^{q}}. \] (20)
For energy influx, we have $T^0_4 < 0$, $T > 0$. We assume implicitly that the form of $T^0_4$ originates mainly in the bulk dynamics, determined through the bulk equations of motion. The presence of the brane generates only a small perturbation to the bulk evolution. This is not expected to be always the case. For example, if the energy density on the brane exceeds a certain value one would expect the influx to stop. This means that $T^0_4$ could vanish or even change sign at the location of the brane, even though it could remain unaffected far from the brane. The complicated dynamics associated with such phenomena is beyond the scope of this work. For our purposes we shall rely on the simple ansatz of eq. (20) that accounts mainly for the dilution of energy density through expansion. The variable $T$ in eq. (20) is taken to be independent of the brane parameters, an approximation that is valid only for finite ranges of the brane evolution.

In this work we do not specify the mechanism that is responsible for the transfer of energy onto the brane. However, the energy influx seems a natural phenomenon in models in which the brane particles are identified with light modes of the five-dimensional theory that are localized on some defect [10]. In a fluctuating system, in which the light modes are not significantly populated, one would expect energy to be transferred from the massive modes to the light ones through interactions. As the light modes are localized on the defect, energy is expected to flow towards it.

The substitution of eq. (20) into eqs. (15)–(17) results in a closed system of equations. In the following we neglect the spatial curvature and set $k = 0$ in eq. (16).

Solutions: We are interested in solutions that describe accelerating eras in the cosmological evolution of the brane, without the presence of an effective cosmological constant. For this reason, we concentrate on the range $0 \leq w \leq 1/3$ for the parameter that determines the equation of state of the brane matter. We also consider the range $0 \leq q \leq 4$ for the parameter appearing in eq. (20), which is directly related to the equation of state of the bulk matter as
we explained above. In a stationary state, in which particles with constant velocity $v_4$ along the fifth dimension fall onto the brane, we expect $q = 3$, $15/4$ or $4$. Slow variations of $v_4$ can be modelled by allowing non-integer values of $q$. For example an increase of $|v_4|$ for infalling massive matter should lead to an increasing parameter $T$ in eq. (20) or, alternatively, to a value of $q$ smaller than 3. We do not specify the details of the mechanism through which energy is transferred from the bulk to the brane. This can be done only within a scenario that includes a dynamical localization mechanism for the brane matter.

For the range $0 \leq w < 1/3$, $0 \leq q \leq 4$ the system of eqs. (15)–(17), (20) has solutions of the form

$$\frac{\rho(a_o)}{6M_{Pl}^2} = \frac{C_1}{a_o^s}, \quad \chi(a_o) = \frac{C_2}{a_o^s}, \quad (21)$$

with

$$s = \frac{2q}{3}, \quad C_1 = \frac{4-s}{3(1+w)-s}, \quad C_2 = \frac{1-3w}{[3(1+w)-s](4-s)T}. \quad (23)$$

Moreover, it can be checked that these solutions are attractors of neighbouring cosmological flows. As a result, the brane evolution is always given by the above equations after a sufficiently long time, independently of the initial conditions. A characteristic property of the solutions is that $C_2 < 0$ and, therefore, $\chi < 0$. On the other hand, $C_1 + C_2 > 0$ and the Hubble parameter is always real.

The solutions (21) describe an expanding Universe with an acceleration parameter

$$Q_o = \frac{1}{H_o^2 a_o} = 1 - \frac{q}{3}. \quad (25)$$

This expression is independent of the equation of state of the brane matter (as long as $w \neq 1/3$). The Hubble parameter, however, depends on $w$, as can be seen from eqs. (16), (24). For $q < 3$ the parameter $Q_o$ is positive and the expansion is accelerating.

For $0 < q \leq 4$ the scale factor increases as a power of $t$ for long times

$$a_o = \left[ a_{oi}^{q/3} + \frac{q}{3} (C_1 + C_2)^{1/2} (t - t_i) \right]^{3/q}. \quad (26)$$

For $q = 0$ the expansion is exponential

$$a_o = a_{oi} e^{(C_1+C_2)^{1/2}(t-t_i)}, \quad (27)$$

independently of the equation of state of the brane matter. This case corresponds to the fixed-point solutions studied in a more general context in ref. [8]. It is remarkable that one can obtain exponential expansion even for non-relativistic brane matter. However, the value $q = 0$ implies the presence of vacuum energy density in the bulk, which is assumed to be transferred to the brane particles.

The case of relativistic brane matter must be considered separately, as eqs. (16), (21), (24) give $H_o \to 0$ for $w \to 1/3$. The solution (21)–(24) is approached only after a very long time in
this limit and the initial conditions determine most of the evolution. For \( w = 1/3 \) the solution is

\[
\frac{1}{6M_{Pl}^2} \rho(a_o) = \frac{1}{6M_{Pl}^2} \rho_{i} \chi_{o}^{4} \frac{T_{o}^{0}}{a_{o}^{2}} + \frac{1}{H_{i} a_{o}^{2} 6 - q} a_{o}^{2-q} \tag{28}
\]

\[
\chi(a_o) = \frac{\chi_{i} \chi_{o}^{4}}{a_{o}^{4}} - \frac{T_{o}^{0}}{H_{i} a_{o}^{2} 6 - q} a_{o}^{2-q}, \tag{29}
\]

where \( a_{oi} = a_{o}(t_{i}) \) etc. The sum \( \rho / (6M_{Pl}^2) + \chi \) that determines the Hubble parameter behaves as an effective radiation term \( \sim a_{o}^{-4} \). This is achieved through the cancellation of the dominant second terms in the r.h.s. of eqs. (28) and (29). As a result, the Universe is always decelerating for \( w = 1/3 \).

**Discussion:** Our main result is summarized by eq. (25): The influx of energy onto the 3-brane from the bulk can lead to accelerated expansion, depending on the equation of state of the bulk matter. For positive acceleration one needs \( q < 3 \) for the variable \( q \) that parametrizes the dependence of the energy flow on the scale factor (see eq. (20)). If the bulk is populated by a gas of non-relativistic particles that drift slowly towards the brane, we expect \( q = 3 \). If the drift velocity increases with time, \( q < 3 \). A value of \( q \) slightly below 3 leads to a positive acceleration parameter smaller than 1, the case favoured by the astronomical data.

The construction of a complete model that realizes this scenario requires the technically difficult solution of the Einstein equations in the bulk. However, certain properties of such a solution can be inferred in general terms. The energy flow along the fifth dimension towards the brane implies the presence of a region of high energy density where the flow originates. For example, one could consider a second brane that emits energy into the bulk. Its evolution is governed by eqs. (15)–(17) with a positive \( T_{0}^{0} \). A solution analogous to that of eqs. (21)–(24) exists, with \( H < 0 \). It predicts a Big Crunch after a finite time. However, neighbouring cosmological flows diverge from this solution. There are also various solutions with \( H > 0 \). The generic cosmological evolution leads to the depletion of the energy density on the second brane. The flow of energy along the fifth dimension is expected to stop after a certain period, whose length depends on the initial energy density of the brane.

These examples indicate that, in a complete model, the solutions (21)–(24) are expected to be valid only for a finite time interval. This makes the derivation of exact analytical expressions more difficult. However, we believe that the physical arguments for the form of the energy-momentum tensor that we gave in the previous section (see eq. (20)) give a valid approximation, especially in the case of slow accretion of energy by the brane.

The RS model is employed here only in order to provide the context in which our arguments can be implemented. The discussion of cosmological solutions must take into account the gravitational effects of the brane tension (the vacuum energy associated with the brane) and the matter localized on it. The RS model is the only example in which this has been realized in such a way that conventional cosmology is reproduced in the low density limit. However, we view this model only as a toy one. Many of its ingredients, such as the negative energy density in the bulk and the structure of the resulting AdS space, do not seem crucial for the emergence of accelerating solutions. This view is supported by the discussion of eq. (1) in the introduction. In order to test this conclusion, it would be interesting to implement this mechanism in a model with a bulk geometry that is flat, at least asymptotically.
An important property of the realization of the mechanism in the context of the RS model is the absence of acceleration during radiation domination, even for energy influx. For brane matter with $w = 1/3$ the solution (28), (29) describes a radiation dominated brane, whose expansion is decelerating. This means that the cosmic acceleration is a phenomenon only of the era of matter domination.

It must be pointed out that the solution (28), (29) depends crucially on the presence of the component $\chi$ of eq. (12). In the case of the pure AdS bulk space of the RS model this component describes, according to the AdS/CFT correspondence, the boundary conformal field theory [11]. It has been characterized as “mirage” or “Weyl radiation”. In our scenario we have assumed a more general matter content for the bulk theory, for which this interpretation is not established. However, the term $\sim 4H\chi$ in the evolution equation of $\chi$ is characteristic of radiation. It is not clear if a similar “radiation” component will be present in a different realization, with an asymptotically flat bulk space for example. Therefore, the absence of acceleration during the radiation dominated era on the brane, even for energy influx, may be a property of RS-type models only.

The energy influx may not be present at all times for other reasons. For example, it may be significant only for very low energy densities on the brane, and be replaced by energy outflow for densities exceeding a critical value. This again would lead to the conclusion that the accelerated expansion is only a recent phenomenon in cosmic terms. The dynamical localization mechanism, that determines the details of energy accretion, is the main direction of further research in order to understand this point and construct phenomenologically viable models.

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