Published by Institute of Physics Publishing for SISSA/ISAS
Received: *February 17, 2005*
Revised: *March 1, 2005*
Accepted: *March 1, 2005*

Linear cosmological perturbations in D-brane gases

Ali Kaya

Department of Physics, Boğaziçi University
34342, Bebek, Istanbul, Turkey
E-mail: ali.kaya@boun.edu.tr

Abstract: We consider linear cosmological perturbations on the background of a D-brane gas in which the compact dimensions and the dilaton are stabilized. We focus on long wavelength fluctuations and find that there are no instabilities. In particular, the perturbation of the internal space performs damped oscillations and decays in time. Therefore, the stabilization mechanism based on D-brane gases in string theory remains valid in the presence of linearized inhomogeneities.

Keywords: D-branes, Classical Theories of Gravity.

Contents

1. Introduction

It is well known that the theory of big-bang nucleosynthesis together with the observational abundances of light elements provide a limit to the change in the size of internal space following the production of these elements. Therefore, string theory should be able to explain why the extra dimensions are stabilized following nucleosynthesis? Intuitively, winding branes are expected to resist the expansion and this may offer an explanation. An accommodation of this idea in early string theory can be found in the paper of Brandenberger and Vafa [1], and recently the stabilization of toroidal extra dimensions by winding strings is verified by explicit calculations in [2, 3] (see also [4]). The model of [1] is developed in [5] to include other extended objects in string theory and this scenario is now known as brane gas cosmology (BGC) (see e.g. [6]–[31] for recent work).

It was also claimed in [1] that the dimensionality of space-time can be explained by the annihilation of winding strings in a three dimensional subspace leading decompactification. From the naive intersection probability of hypersurfaces embedded in d-spatial dimensions one can infer that $d = 2p + 1$ is the critical value which allows annihilation of p-branes. Remarkably, for strings $d = 3$ and according to [1] that is why a three dimensional subspace is decompactified. As an additional consequence of this argument, in BGC all higher dimensional branes are expected to be annihilated in the early universe. Although these claims are supported numerically in [32], there are also some caveats. It is observed in [15] that the number of decompactified dimensions depends on initial conditions and the outcome of three large dimensions is argued to be statistically insignificant [25]. In [26] it is shown that the interaction rates of strings are not strong enough to sustain thermal equilibrium. Moreover, D-brane anti-D-brane annihilation rate turns out to be small compared to the expansion rate of a FRW universe including the higher dimensional D-branes [12]. Therefore, it seems that the dimensional counting argument of [1] is not sufficient alone for the annihilation process and it is plausible to have surviving branes other than strings or membranes.

In the light of these results, we think that late time models with higher dimensional branes should also be studied in the context of BGC. Such models may be useful in solving two difficulties with string gases. Firstly, as pointed out in [7], strings may not be able to stabilize volume modulus, or further, there may not be any winding strings in the spectrum depending on the topology of internal space (consider e.g. a sphere). Secondly, it turns out to be difficult to stabilize the dilaton and keep up the stability of extra dimensions [23], which may cause problems in canonical Einstein frame after compactification.

In [24], we have considered a gas of D6-branes wrapping over the six-dimensional internal space and demonstrated dynamical stabilization of the volume modulus and the dilaton. The aim of this paper is to consider long wavelength perturbations in this model to see the validity of the conclusions under linearized inhomogeneities. For the string gas the perturbation analysis has been carried out in [18, 19]. The plan of the paper is as follows. In the next section we derive linearized equations for the scalar metric and dilaton perturbations in the generalized longitudinal gauge. In section 3, we focus on long wavelength fluctuations and solve the equations. We briefly conclude in section 4.

2. Perturbation equations

We start with ten-dimensional dilaton gravity coupled to matter Lagrangian $\mathcal{L}_m$ representing brane winding and momentum modes:

$$S = \int d^{10}x \sqrt{-g}\, e^{-2\phi} \left[R + 4(\nabla\phi)^2 + e^{a\phi}\mathcal{L}_m\right]. \tag{2.1}$$

The field equations that follow from the above action can be found as

$$R_{\mu\nu} + 2\nabla_\mu\nabla_\nu\phi - \frac{1}{2}g_{\mu\nu}\left[R + 4\nabla^2\phi - 4(\nabla\phi)^2\right] = e^{a\phi}T_{\mu\nu}\,, \tag{2.2}$$

$$R + 4\nabla^2\phi - 4(\nabla\phi)^2 = \frac{(a-2)}{2}\, e^{\phi}\mathcal{L}_m\,, \tag{2.3}$$

$$\nabla^\mu T_{\mu\nu} = \frac{(a-2)}{2}\mathcal{L}_m\nabla_\nu\phi - (a-2)T_{\nu\lambda}\nabla^\lambda\phi\,, \tag{2.4}$$

where $T_{\mu\nu}$ is the matter energy-momentum tensor:

$$T_{\mu\nu} = -\frac{1}{\sqrt{-g}}\frac{\partial}{\partial g^{\mu\nu}}\left(\sqrt{-g}\,\mathcal{L}_m\right). \tag{2.5}$$

Equations (2.2) and (2.3) are obtained by directly varying the action with respect to the metric and the dilaton. Although for matter equations one in principle needs Lagrangian $\mathcal{L}_m$ explicitly, in our case (2.4) can be deduced from the consistency of (2.2) and (2.3).

It is well known that one particular combination of the field equations gives a constraint on initial data. Let n^μ to denote the unit normal vector field to an initial value hypersurface. One can easily show that (2.2) contracted with $n^\mu n^\nu$

$$\left(R_{\mu\nu} + 2\nabla_\mu\nabla_\nu\phi - \frac{1}{2}g_{\mu\nu}\left[R + 4\nabla^2\phi - 4(\nabla\phi)^2\right] - e^{a\phi}\,T_{\mu\nu}\right)n^\mu n^\nu = 0 \tag{2.6}$$

does not contain any second order time derivatives. Firstly, it is well from the initial value formulation of general relativity that $n^\mu n^\nu(R_{\mu\nu} - g_{\mu\nu}R/2)$ involves only first order time derivatives of the metric (see e.g. [33]). On the other hand, the dangerous terms with two derivatives acting on dilaton in (2.6) can be reexpressed straightforwardly as

$$\begin{aligned}\nabla^2\phi + n^\mu n^\nu \nabla_\mu\nabla_\nu\phi &= h^{\mu\nu}\nabla_\mu\nabla_\nu\phi\,,\\ &= D^\mu D_\mu\phi - (D_\mu n^\mu)(n^\nu\nabla_\nu\phi)\,,\end{aligned} \tag{2.7}$$

where $g_{\mu\nu} = h_{\mu\nu} - n_\mu n_\nu$ and $D_\mu = h_\mu{}^\nu\nabla_\nu$ is the spatial covariant derivative along constant time hypersurface. It is clear that (2.7) contains only first order time derivatives and thus (2.6) can be viewed as a constraint on initial data.

Assuming a metric of the following form

$$ds^2 = -e^{2A}dt^2 + e^{2B}dx^i dx^i + e^{2C}\,d\Sigma_p^2\,, \tag{2.8}$$

$(i = 1, \ldots, m)$ the energy-momentum tensor for a gas of branes winding over the compact manifold Σ_p can be found as (see e.g. [22])

$$\begin{aligned}T_{\hat{0}\hat{0}} &= T_w\,e^{-mB} + T_m e^{-mB-(p+1)C}\,,\\ T_{\hat{i}\hat{j}} &= 0\,,\\ T_{\hat{a}\hat{b}} &= -T_w e^{-mB}\,\delta_{ab} + \frac{T_m}{p}\,e^{-mB-(p+1)C}\,\delta_{ab}\,,\end{aligned} \tag{2.9}$$

where T_w and T_m are constants corresponding to winding and momentum modes. In this paper we take Σ_p to be Ricci flat. From (2.9) an equation of state $p_{\hat{\mu}} = \omega_\mu\rho$ can be deduced where $\omega_i = 0, \omega_a = -1$ for winding modes and $\omega_i = 0, \omega_a = 1/p$ for momentum modes. Note that the equation for momentum modes is identical to radiation confined in the compact space. One can check from (2.5) and (2.9) that

$$\mathcal{L}_m = -2\rho\,, \tag{2.10}$$

which is precisely the Lagrangian for hydrodynamical matter (see e.g. section 10.2 of [34]). In general we have $\nabla^\mu T_{\mu\nu} \neq 0$ due to the coupling of $\mathcal{L}_m$ to dilaton.

To focus on D-branes one should set

$$a = 1, \tag{2.11}$$

which is required to have the correct $g_s = e^\phi$ dependence in the low energy effective action (2.1). Recall that we also have

$$m + p = 9, \tag{2.12}$$

where m and p are the dimensions of the observed and the compact dimensions, respectively.

In [24], we have found a particular solution to (2.2)-(2.4)

$$\begin{aligned}ds^2 &= -dt^2 + (t)^{8/(m+3)}\,dx^i dx^i + e^{2C_0}\,d\Sigma_p^2\,,\\ e^\phi &= e^{\phi_0}\,(t)^{2(m-3)/(m+3)}\,,\\ e^{\phi_0} &= \frac{4\,(p+2)\,p}{T_w\,(p+1)(m+3)^2}\,, \qquad e^{(p+1)C_0} = \frac{(p+2)\,T_m}{p\,T_w}\,,\end{aligned} \tag{2.13}$$

which has some appealing properties. For instance, it is possible to argue analytically and support the argument numerically that for arbitrary initial conditions (within a given ansatz) the fields approach (2.13) asymptotically. Moreover, for $m = 3$, the dilaton becomes a constant and thus the canonical Einstein and string frames become identical up to a constant scaling.

Let us now consider linearized inhomogeneities on (2.13). We are interested in the case where the fluctuations are independent of compact coordinates[1] and choose the generalized longitudinal gauge (see [18]) in which the scalar metric perturbations only appear in diagonal entries. Therefore, the perturbed solution has the form:

$$\begin{aligned} ds^2 &= -e^{2A}(1+2\delta A)dt^2 + e^{2B}(1+2\delta B)\,dx^i dx^i + e^{2C}(1+2\delta C)d\Sigma_p^2\,, \\ e^{\Phi} &= e^{\phi}(1+\delta\phi)\,, \end{aligned} \tag{2.14}$$

where A, B, C and ϕ denote the background values depending only on t and

$$\delta A = \delta A(t,x^i)\,, \qquad \delta B = \delta B(t,x^i)\,, \qquad \delta C = \delta C(t,x^i)\,, \qquad \delta\phi = \delta\phi(t,x^i)\,. \tag{2.15}$$

Let us remind that in the longitudinal gauge δA, δB and δC coincide with diffeomorphism invariant quantities.

In obtaining the linearized field equations we find it convenient to use (2.6) together with a particular combination of (2.2) and (2.3)

$$R_{\mu\nu} + 2\nabla_\mu\nabla_\nu\phi = \frac{1}{2}g_{\mu\nu}e^{\phi}\rho + e^{\phi}\,T_{\mu\nu}. \tag{2.16}$$

As we pointed out above, (2.6) can be interpreted as a constraint on initial data.[2] Substituting (2.9) and (2.14) in these equations and keeping only the linear terms one gets

$$\begin{aligned} &2\ddot{\delta\phi} - m\,\ddot{\delta B} - p\,\ddot{\delta C} + e^{-2B}\sum_k \delta A_{kk} + 2\left[m\ddot{B} + m\,\dot{B}\dot{B} - 2\ddot{\phi}\right]\delta A + \\ &\qquad\qquad + \left[m\dot{B} - 2\dot{\phi}\right]\dot{\delta A} - 2m\dot{B}\dot{\delta B} = \frac{1}{2}\rho e^{\phi}\delta\phi + \frac{1}{2}e^{\phi}\delta\rho\,, \end{aligned} \tag{2.17}$$

$$(m-1)\dot{\delta B}_i - 2\dot{\delta\phi}_i - \left[(m-1)\dot{B} - 2\dot{\phi}\right]\delta A_i + p\dot{\delta C}_i - p\dot{B}\delta C_i + 2\dot{B}\delta\phi_i = -e^{\phi+B}\delta T_{0i}\,, \tag{2.18}$$

$$\begin{aligned} &\ddot{\delta B} + e^{-2B}\left[2\delta\phi_{ii} - \delta A_{ii} - p\delta C_{ii} - (m-2)\delta B_{ii}\right] - e^{-2B}\sum_k \delta B_{kk} + 2\left[m\dot{B} - \dot{\phi}\right]\dot{\delta B} - \\ &\qquad\qquad - \dot{B}\dot{\delta A} + p\dot{B}\dot{\delta C} - 2\dot{B}\dot{\delta\phi} + 2\left[2\dot{B}\dot{\phi} - \ddot{B} - m\dot{B}\dot{B}\right]\delta A = \\ &\qquad\qquad\qquad\qquad = \frac{1}{2}\rho e^{\phi}\delta\phi + \frac{1}{2}e^{\phi}\delta\rho\,, \end{aligned} \tag{2.19}$$

$$\left[2\delta\phi - \delta A - p\delta C - (m-2)\delta B\right]_{ij} = 0, \qquad i \neq j \tag{2.20}$$

$$\ddot{\delta C} - e^{-2B}\sum_k \delta C_{kk} + \left[m\dot{B} - 2\dot{\phi}\right]\dot{\delta C} = \frac{1}{2}\rho e^{\phi}\delta\phi + \frac{1}{2}e^{\phi}\delta\rho + e^{\phi}T_{aa}\delta\phi + e^{\phi}\delta T_{aa}\,, \tag{2.21}$$

---

[1] This is indeed a good approximation since the extra dimensions are assumed to be very small.

[2] It is possible to show that (2.6) and (2.16) are equivalent to (2.2) and (2.3) with $a = 1$.

$$\left[m(m-1)\dot{B}-2m\dot{\phi}\right]\dot{\delta B}+\left[mp\dot{B}-2p\dot{\phi}\right]\dot{\delta C}+\left[4m\dot{B}\dot{\phi}-4\dot{\phi}\dot{\phi}-m(m-1)\dot{B}\dot{B}\right]\delta A+$$
$$+\left[4\dot{\phi}-2m\dot{B}\right]\dot{\delta\phi}+e^{-2B}\sum_k\left[2\delta\phi_{kk}-(m-1)\delta B_{kk}-p\delta C_{kk}\right]=$$
$$=\rho e^{\phi}\delta\phi+e^{\phi}\delta\rho\,, \qquad (2.22)$$

where (2.17)–(2.21) follow from (2.16), and (2.22) can be obtained from (2.6). In the above equations the dot and the subindex i denote partial differentiations with respect to t and x^i, respectively, and we have not used the summation convention.

Note that in (2.18) we introduce a non-diagonal component for the linearized energy momentum tensor δT_{0i}. Although for long wavelength fluctuations this equation is satisfied identically, it is in general necessary to add δT_{0i} since the perturbed metric functions depend on both t and x^i, which gives non-zero R_{0i} (indeed (2.18) is the $(0i)$ component of (2.16)). Moreover, one can see that without δT_{0i} the matter equations constrain the metric functions too much, see (2.24) below. The appearance of this non-diagonal term for winding modes can be deduced from the Born-Infeld (BI) action as follows. When the metric functions depend only on t, a static brane winding Σ_p and located at a constant position $x^i=x^i_0$ becomes an extremum of the BI action (i.e. the embedding coordinates become harmonic maps). Using this solution in the energy-momentum tensor obtained from BI action and smearing the delta function singularity by taking a continuum average for a gas of branes, one gets the winding part of (2.9) [9]. However, when the metric is perturbed with t and x^i dependent functions, the static brane solution should be modified $x^i=x^i_0+\delta x^i(t)$ and this correction induces non-zero δT_{0i}.

Linearizing the matter equation (2.4) we find that $\nu=a$ component is identically satisfied while $\nu=0$ and $\nu=i$ equations give

$$\dot{\delta\rho}+(m\dot{B}+p\dot{C})\delta\rho+(m\dot{\delta B}+p\dot{\delta C})\rho+\dot{\delta C}\,T^a{}_a+\dot{C}\,\delta T^a{}_a+e^{A-B}\delta T^{0j}_{,j}=0\,, \qquad (2.23)$$
$$\dot{\delta T}_{0i}+\left[(m+1)\dot{B}+p\dot{C}-\dot{\phi}\right]\delta T_{0i}+e^{A-B}\left[\delta C_i T^a{}_a+\rho\delta\phi_i-\rho\delta A_i\right]=0\,. \qquad (2.24)$$

Equation (2.23) has actually two separate parts corresponding to winding and momentum modes. In obtaining $\delta\rho$ and δT_{ab} from (2.9), one should also consider the variations δT_w and δT_m. Therefore, there are seven variables which are δT_w, δT_m, δT_{0i}, δA, δB, δC and $\delta\phi$. Equations (2.23) and (2.24) determine the evolution of source fields δT_w, δT_m and δT_{0i}, which in turn should be substituted in (2.17)–(2.22). It is possible to solve (2.20) as

$$2\delta\phi-\delta A-p\,\delta C-(m-2)\delta B=0\,, \qquad (2.25)$$

which reduces the number of unknown functions by one. On the other hand, (2.18) and (2.22) are not dynamical evolution equations since they only contain first order time derivatives and thus merely restrict the solution space. Ignoring them, it is easy to see that the number of equations and unknowns are the same.

3. Long wavelength fluctuations

Let us now focus on long wavelength perturbations. As pointed out in [18] an instability on such scales would be a problem for the stabilization mechanism. In this regime all spatial derivatives can be ignored in field equations which in turn implies that (2.18) and (2.24) are identically satisfied. On the other hand, from (2.23) one obtains

$$\delta\dot{T}_w = \delta\dot{T}_m = 0\,. \tag{3.1}$$

Since constant shifts of T_w or T_m alter the solution (2.13), they are not honest perturbations and one should set $\delta T_w = \delta T_m = 0$. In other words, (3.1) gives the zero modes which must be ignored. Solving $\delta\phi$ from (2.25)[3] and using the background values from (2.13), we find that (2.17), (2.19), (2.21) and (2.22) respectively become

$$\delta\ddot{A} - 2\delta\ddot{B} + \frac{12\delta\dot{A}}{(m+3)t} - \frac{8m\delta\dot{B}}{(m+3)t} - \frac{10p\delta A}{(m+3)^2t^2} + \frac{2p(m+2)\delta B}{(m+3)^2t^2} = 0\,, \tag{3.2}$$

$$\delta\ddot{B} - \frac{8\delta\dot{A}}{(m+3)t} + \frac{20\delta\dot{B}}{(m+3)t} - \frac{10p\delta A}{(m+3)^2t^2} + \frac{2p(m+2)\delta B}{(m+3)^2t^2} = 0\,, \tag{3.3}$$

$$\delta\ddot{C} + \frac{12\delta\dot{C}}{(m+3)t} + \frac{2p(p+2)\delta C}{(m+3)^2t^2} = 0\,, \tag{3.4}$$

$$\frac{12\delta\dot{A}}{(m+3)t} + \frac{4(m-6)\delta\dot{B}}{(m+3)t} + \frac{20p\delta A}{(m+3)^2t^2} - \frac{4p(m+2)\delta B}{(m+3)^2t^2} = 0\,. \tag{3.5}$$

Note that δC decouples from δA and δB. As a consistency check of these equations, we find that the time derivative of the constraint (3.5) is proportional to four times (3.2) plus m times (3.3) and thus identically satisfied, as it should be. This does not mean however that (3.5) is empty. Rather, one should first solve (3.2)-(3.3) and then use (3.5) to narrow the solution space.

Equation (3.4) can easily be solved for δC which gives

$$\delta C(t) = c_1\,\frac{\cos(a\ln t)}{t^b} + c_2\,\frac{\sin(a\ln t)}{t^b}\,, \tag{3.6}$$

where a, b are positive numbers depending on m (e.g. for $m = 3$, $a = \sqrt{29/12}$, $b = 1/2$). On the other hand, the coupled equations (3.2) and (3.3) give

$$\begin{aligned} \delta A &= (m+2)\,k_1 + k_2\,t^{-(7+m)/(m+3)} + k_3(m-2)t^{(-9+m)/(m+3)} + k_4\,(29+11m)t^2 \\ \delta B &= 5k_1 - k_2 t^{-(7+m)/(m+3)} + k_3\,t^{(-9+m)/(m+3)} + k_4(23+m)t^2 \end{aligned} \tag{3.7}$$

where $k_1, \ldots, k_4$ are arbitrary constants. Substituting these solutions into (3.5), we observe that it yields

$$k_4 = 0\,, \tag{3.8}$$

and thus the growing mode in (3.7) is eliminated by the constraint equation. As a result, we find that there are no instabilities to this order.

[3]For long wavelength fluctuations (2.25) can be viewed as the residual gauge symmetry fixing.

4. Conclusions

In this paper, we consider long wavelength fluctuations on the background of a D-brane gas in ten dimensional dilaton gravity constructed in [24]. The most attractive feature of the solution is that it yields stabilized internal dimensions. Moreover, when the observed space is three dimensional the dilaton also becomes a constant. The main result of this paper is that the long wavelength perturbation of the internal space performs damped oscillations. Other fluctuations, including the dilaton, are found to decay to a constant. Therefore, the stabilization mechanism based on D-brane winding and momentum modes remains valid in the presence of linearized inhomogeneities.

Technically, the perturbation analysis for D-branes is similar to the string gases studied in [18]. The main difference is that in our case there appears to be a growing mode which is eliminated by the constraint equation. Although the presence of this mode would not affect the internal space, it would ruin the stabilization of dilaton for $m = 3$, which is crucial in the effective field theory obtained after compactification.

Higher dimensional D-brane gases offer an analogous scenario to string gases for the stabilization of the extra dimensions. Moreover, D6 branes happen to fix dilaton in three dimensions. This mechanism can be useful in topologically non-trivial compactifications where the winding strings do not appear in the spectrum or they are not capable of fixing volume modulus. Of course, the toy model we consider is immature and it should be developed in many different directions. However, the results in BGC show that string theory has the right ingredients to solve the stabilization problem of extra dimensions in cosmology.

References

[1] R. Branderberger and C. Vafa, *Superstrings in the early universe*, *Nucl. Phys.* **B 316** (1989) 391.

[2] S. Watson and R. Brandenberger, *Stabilization of extra dimensions at tree level*, *JCAP* **11** (2003) 008 [hep-th/0307044].

[3] S.P. Patil and R. Brandenberger, *Radion stabilization by stringy effects in general relativity and dilaton gravity*, hep-th/0401037.

[4] S.P. Patil and R.H. Brandenberger, *The cosmology of massless string modes*, hep-th/0502069.

[5] S. Alexander, R.H. Brandenberger and D. Easson, *Brane gases in the early universe*, *Phys. Rev.* **D 62** (2000) 103509 [hep-th/0005212].

[6] R. Brandenberger, D.A. Easson and D. Kimberly, *Loitering phase in brane gas cosmology*, *Nucl. Phys.* **B 623** (2002) 421 [hep-th/0109165].

[7] D.A. Easson, *Brane gases on K3 and Calabi-Yau manifolds*, *Int. J. Mod. Phys.* **A 18** (2003) 4295 [hep-th/0110225].

[8] D.A. Easson, *Brane gas cosmology and loitering*, hep-th/0111055.

[9] A. Kaya and T. Rador, *Wrapped branes and compact extra dimensions in cosmology*, *Phys. Lett.* **B 565** (2003) 19 [hep-th/0301031].

JHEP03(2005)003

[10] B.A. Bassett, M. Borunda, M. Serone and S. Tsujikawa, *Aspects of string-gas cosmology at finite temperature*, *Phys. Rev.* **D 67** (2003) 123506 [hep-th/0301180].

[11] A. Kaya, *On winding branes and cosmological evolution of extra dimensions in string theory*, *Class. and Quant. Grav.* **20** (2003) 4533 [hep-th/0302118].

[12] M. Majumdar and A.-C. Davis, *D-brane anti-brane annihilation in an expanding universe*, *J. High Energy Phys.* **12** (2003) 012 [hep-th/0304153].

[13] A. Campos, *Late-time dynamics of brane gas cosmology*, *Phys. Rev.* **D 68** (2003) 104017 [hep-th/0304216].

[14] R. Brandenberger, D.A. Easson and A. Mazumdar, *Inflation and brane gases*, *Phys. Rev.* **D 69** (2004) 083502 [hep-th/0307043].

[15] R. Easther, B.R. Greene, M.G. Jackson and D. Kabat, *Brane gases in the early universe: thermodynamics and cosmology*, *JCAP* **01** (2004) 006 [hep-th/0307233].

[16] T. Biswas, *Cosmology with branes wrapping curved internal manifolds*, *J. High Energy Phys.* **02** (2004) 039 [hep-th/0311076].

[17] A. Campos, *Late cosmology of brane gases with a two-form field*, *Phys. Lett.* **B 586** (2004) 133 [hep-th/0311144].

[18] S. Watson and R. Brandenberger, *Linear perturbations in brane gas cosmology*, *J. High Energy Phys.* **03** (2004) 045 [hep-th/0312097].

[19] S. Watson, *Uv perturbations in brane gas cosmology*, *Phys. Rev.* **D 70** (2004) 023516 [hep-th/0402015].

[20] T. Battefeld and S. Watson, *Effective field theory approach to string gas cosmology*, *JCAP* **06** (2004) 001 [hep-th/0403075].

[21] J.Y. Kim, *Late time evolution of brane gas cosmology and compact internal dimensions*, *Phys. Rev.* **D 70** (2004) 104024 [hep-th/0403096].

[22] A. Kaya, *Volume stabilization and acceleration in brane gas cosmology*, *JCAP* **08** (2004) 014 [hep-th/0405099].

[23] A.J. Berndsen and J.M. Cline, *Dilaton stabilization in brane gas cosmology*, hep-th/0408185.

[24] S. Arapoglu and A. Kaya, *D-brane gases and stabilization of extra dimensions in dilaton gravity*, *Phys. Lett.* **B 603** (2004) 107 [hep-th/0409094].

[25] R. Easther, B.R. Greene, M.G. Jackson and D. Kabat, *String windings in the early universe*, *JCAP* **02** (2005) 009 [hep-th/0409121].

[26] R. Danos, A.R. Frey and A. Mazumdar, *Interaction rates in string gas cosmology*, *Phys. Rev.* **D 70** (2004) 106010 [hep-th/0409162].

[27] R. Brandenberger, Y.-K.E. Cheung and S. Watson, *Moduli stabilization with string gases and fluxes*, hep-th/0501032.

[28] A. Campos, *Asymptotic cosmological solutions for string/brane gases with solitonic fluxes*, hep-th/0501092.

[29] T. Biswas, R. Brandenberger, D.A. Easson and A. Mazumdar, *Coupled inflation and brane gases*, hep-th/0501194.

[30] T. Rador, *Intersection democracy for winding branes and stabilization of extra dimensions*, hep-th/0501249.

[31] T. Rador, *Vibrating winding branes, wrapping democracy and stabilization of extra dimensions in dilaton gravity*, hep-th/0502039.

[32] M. Sakellariadou, *Numerical experiments in string cosmology*, *Nucl. Phys.* **B 468** (1996) 319 [hep-th/9511075].

[33] R.M. Wald, *General relativity*, The University of Chicago Press, Chicago 1984.

[34] V.F. Mukhanov, H.A. Feldman and R.H. Brandenberger, *Theory of cosmological perturbations. Part 1, classical perturbations. Part 2, quantum theory of perturbations. Part 3, extensions*, *Phys. Rept.* **215** (1992) 203.