Abstract

Unified dark matter/energy models (quartessence) based upon the Chaplygin gas D-brane fail owing to the suppression of structure formation by the adiabatic speed of sound. Including string theory effects, in particular the Kalb-Ramond field which becomes massive via the brane, we show how nonadiabatic perturbations allow successful structure formation.

1 Introduction

Cosmology today faces a mystery, dark matter, wrapped in an enigma: dark energy. The existence of dark matter has been known for many years from galactic rotation curves and cluster dynamics, with the lightest supersymmetric particle being the leading suspect. Dark energy is a comparatively recent discovery coming from observations of high-redshift supernovae [1], and is held responsible for the current accelerated Hubble expansion. Canonical models of dark energy include a cosmological constant, quintessence, and k-essence; for recent reviews, see [2]. Tight constraints come from the WMAP [3] observations of the cosmic microwave background (CMB): $\Omega_{\text{DE}} = 0.72$, $w_{\text{DE}} = p_{\text{DE}}/\rho_{\text{DE}} < -0.78$, $\Omega_{\text{DM}} = 0.24$, $w_{\text{DM}} \approx 0$, with an impurity of 4% of ordinary baryonic matter.

In contrast to the standard assumption that dark matter and dark energy are distinct, there stands the hypothesis that both are different manifestations of a single entity. The first definite model of this type [4] is based on the Chaplygin gas, an exotic fluid with an equation of state $p = -A/\rho$, whose cosmological possibilities were earlier noticed by Kamenshchik et al. [5]. Subsequently, the generalization to $p = -A/\rho^\alpha$, $0 \leq \alpha \leq 1$, was given [6] and the term ‘quartessence’ coined [7] to describe unified dark matter/energy models.

In a homogeneous model [6], the Chaplygin gas density is $\rho(a) = \sqrt{A + B/a^6}$, where $B$ is an integration constant and $a$ the scale factor normalized to unity today, thus interpolating between dust ($\rho \sim a^{-3}$) and a cosmological constant ($\rho \sim \sqrt{A}$). The inhomogeneous...
Chaplygin gas based on a Zel’dovich type approximation has been proposed \[4\], and the picture has emerged that on caustics, where the density is high, the fluid behaves as cold dark matter, whereas in voids, \( w = p/\rho \) is driven to the lower bound \(-1\) producing acceleration as dark energy. Soon, however, it has been shown that the simple Chaplygin gas model does not work \[8, 9, 10\]. The physical reason is that while the adiabatic speed of sound, defined by \( c^2_s = \frac{\partial p}{\partial \rho}\big|_S \), is small until \( a \sim 1 \), the accumulated comoving acoustic horizon \( d_s = \int \frac{dt c_s}{a} \approx H_0^{-1} a^{7/2} \) reaches Mpc scales by redshifts of twenty, frustrating structure formation even into the mildly nonlinear regime \[11\]. In the absence of caustic formation, the simple Chaplygin gas undergoes damped oscillations that are in gross conflict with the observed mass power spectrum \[8\] and the CMB \[9, 10\]. The problem is not alleviated by the generalized Chaplygin gas, \( d_s \approx \sqrt{\alpha H_0^{-1} a^{2+3\alpha/2}} \), unless \( \alpha \) is fine-tuned to unnaturally small values.

Such is the appeal of the quintessence idea that several variations have been proposed to overcome the acoustic barrier. In particular, Scherrer \[12\] has explored a \( k \)-essence type model where the Lagrangian has a local minimum as a function of the derivatives of the \( k \)-essence field; such a model is equivalent to the ghost condensate \[13\] and hence shares its peculiarities \[14\]. Barreira and Sen \[15\] have suggested the generalized Chaplygin gas in a modified gravity approach, reminiscent of Cardassian models \[16\]. Yet, another deformation of the Chaplygin gas is described in \[17\]. One simple way to save the Chaplygin gas is to suppose that nonadiabatic perturbations cause the pressure perturbation \( \delta p \) to vanish, and with it the acoustic horizon \[18\]. To achieve this, it is necessary to add new degrees of freedom which, to some extent, spoil the simplicity of quintessence unification.

One of the most appealing aspects of the original Chaplygin gas model is that it is equivalent \[4\] to the Dirac-Born-Infeld description of a D-brane in string theory \[19\]. In fact, string theory branes possess two features that are absent in the simple Chaplygin gas: (i) they support an Abelian gauge field \( A_\mu \) reflecting open strings with their ends stuck on the brane; (i) they couple to the (pull-back of) Kalb-Ramond \[20\] antisymmetric tensor field \( A_{\mu \nu} = -A_{\nu \mu} \) which, like the gravitational field \( g_{\mu \nu} \), belongs to the closed string sector. In this paper we show that when these natural stringy features are incorporated into Chaplygin quintessence, the basic scenario previously proposed in \[4\] is realized owing to nonadiabatic perturbations. It is important to stress that both the Dirac-Born-Infeld and Kalb-Ramond fields originate from an ultimate theory in the context of string/M theory and, unlike in quintessence models, each field affects both dark matter and dark energy. Hence, dark matter and dark energy remain unified as in the simple Chaplygin gas model.

We organize the paper as follows. In Sec. 2 we outline the basic ideas behind the nonadiabatic perturbation mechanism. In Sec. 3 we describe the quartessence model based on the Dirac-Born-Infeld Lagrangian extended by the Kalb-Ramond field. Using the Newtonian approach for simplicity, which is appropriate when the pressure is small compared with the density, we demonstrate the cancellation of the pressure perturbation by the nonadiabatic perturbation mechanism. A discussion is given in Sec. 4 and Sec. 5 concludes the paper.
2 Nonadiabatic perturbations

It has been noted by Reis et al. [18] that the root of the structure formation problem is the term $\Delta \delta p$ in perturbation equations, equal to $c_s^2 \Delta (\delta \rho/\bar{\rho})$ for adiabatic perturbations, and if there are entropy perturbations such that $\delta p = 0$, no difficulty arises. This scenario, which is difficult to justify in the simple Chaplygin gas model, may be realized by introducing an extra degree of freedom, e.g., in terms of a quintessence-type scalar field $\phi$.

Suppose that the matter Lagrangian depends on two degrees of freedom, e.g., a Born-Infeld scalar field $\theta$ and one additional scalar field $\phi$. In this case, instead of a simple barotropic form $p = p(\rho)$, the equation of state involves the entropy density (entropy per particle) $s$ and may be written in the parametric form

$$p = p(\theta, \phi),$$
$$\rho = \rho(\theta, \phi),$$
$$s = s(\theta, \phi).$$

One of the parameters, e.g., $\theta$ may be eliminated in terms of $\rho$, so that the two remaining independent quantities, the pressure $p$ and the entropy density $s$, are expressed as functions of $\rho$ and $\phi$ and, hence, the corresponding perturbations are

$$\delta p = \delta \rho \frac{\partial p}{\partial \rho} + \delta \phi \frac{\partial p}{\partial \phi},$$
$$\delta s = \delta \rho \frac{\partial s}{\partial \rho} + \delta \phi \frac{\partial s}{\partial \phi}.\quad (4,5)$$

Then, the speed of sound squared

$$c_s^2 \equiv \frac{\delta p}{\delta \rho} \bigg|_{\delta s = 0} = \frac{\partial p}{\partial \rho} - \frac{\partial s}{\partial \rho} \left( \frac{\partial s}{\partial \phi} \right)^{-1} \frac{\partial p}{\partial \phi}.\quad (6)$$

is expressed as the sum of two nonadiabatic terms: the usual $\partial p/\partial \rho$ and one additional term due to the nonadiabatic perturbation of the field $\phi$. Thus, even for a nonzero $\partial p/\partial \rho$, the speed of sound may vanish if the second term on the right-hand side of (6) cancels the first one. This cancellation will take place if in the course of an adiabatic expansion, the perturbation $\delta \phi$ grows with $a$ in the same way as $\delta \rho$. In this case, it is only a matter of adjusting initial conditions of $\delta \phi$ with $\delta \rho$ to get $c_s = 0$.

To illustrate the above general discussion, consider the recently proposed extended quintessence model [21] with the Lagrangian

$$\mathcal{L}_{EQ} = \frac{1}{2} g^{\mu\nu} \phi_{,\mu} \phi_{,\nu} - \sqrt{A} e^{-\omega \phi} \sqrt{1 - g^{\mu\nu} \theta_{,\mu} \theta_{,\nu}},$$

where $\omega$ is a coupling constant. That is to say, we have replaced the constant brane tension, or the Chaplygin gas constant, by the potential for a quintessence-type field $\phi$, here taken to
be exponential. The model of Eq. (7) is constructed so that the perturbation of the pressure associated to the field $\theta$ is given by

$$\delta p_\theta = -\bar{p}_\theta \left( \frac{\delta \rho_\theta}{\rho_\theta} - 2\omega \delta \phi \right)$$

(8)

and the above scenario is realized by $\delta \phi = \delta \rho_\theta/(2\omega \bar{\rho}_\theta)$ as an initial condition outside the causal horizon $d_c = \int dt/a \simeq H_0^{-1} a^{1/2}$. However, the perturbation $\delta \phi$ satisfies

$$\ddot{\delta \phi} + 3H \dot{\delta \phi} - \frac{1}{a^2} \Delta \delta \phi \simeq 0,$$

(9)

with the solution in $k$-space

$$\delta \phi_k = a^{-3/4} J_{3/2}(kd_c).$$

(10)

Then, once the perturbations enter the causal horizon $d_c$ (but are still outside the acoustic horizon $d_s$), $\delta \phi$ undergoes rapid damped oscillations, so that the nonadiabatic perturbation associated with $\phi$ is destroyed. This means that the nonadiabatic perturbations are not automatically preserved except on long, i.e., superhorizon, wavelengths where the simple Chaplygin gas has no problem anyway. Thus, the extended quartessence model of Eq. (7) unfortunately does not solve the structure formation problem.

In the following section we demonstrate how the Kalb-Ramond field provides a mechanism for the desired cancellation of the nonadiabatic perturbations.

### 3 New Quartessence

We begin by recalling that the Kalb-Ramond field strength $H_{\mu\nu\alpha} = A_{\mu\nu,\alpha} + A_{\alpha\mu,\nu} + A_{\nu\alpha,\mu}$ is invariant under $A_{\mu\nu} \rightarrow A_{\mu\nu} + \xi_{\mu,\nu} - \xi_{\nu,\mu}$, with a further $\xi_{\mu} \rightarrow \xi_{\mu} + \phi_{,\mu}$ gauge invariance. Thus, in four space-time dimensions there is one (6 - 4 - 1 = 1) degree of freedom, equivalent to a pseudoscalar. The D3 brane is described by

$$S_{DBI} = \int d^4x \sqrt{-\det(g_{\mu\nu})} \mathcal{L}_{DBI} = -\sqrt{A} \int d^4x \sqrt{-\det(g^{(ind)}_{\mu\nu}) + B_{\mu\nu}},$$

(11)

where $g^{(ind)}_{\mu\nu}$ is the induced metric on the brane, $B_{\mu\nu} = A_{\mu\nu} + F_{\mu\nu}, F_{\mu\nu} = A_{\nu,\mu} - A_{\mu,\nu}$ (we have absorbed some factors into $A_{\mu}$). The Dirac-Born-Infeld Lagrangian is invariant under Kalb-Ramond gauge transformations if $A_{\mu} \rightarrow A_{\mu} + \xi_{\mu}$ simultaneously. Of course, the U(1) invariance $A_{\mu} \rightarrow A_{\mu} + \phi_{,\mu}$ gives two degrees of freedom to $A_{\mu}$. Through the Higgs mechanism, $A_{\mu\nu}$ may eliminate $A_{\mu}$ to yield three massive degrees of freedom, as is made manifest by noting that the Bianchi identity for $F_{\mu\nu}$ gives

$$H_{\mu\nu\alpha} = B_{\mu\nu,\alpha} + B_{\alpha\mu,\nu} + B_{\nu\alpha,\mu},$$

(12)

whereas in the temporal (unitary) gauge, $B_{0i} = 0, B_{ij} = \epsilon_{ijk}B^k$. Similar observations have been made in [22].

The matter is then described by the Lagrangian

$$\mathcal{L} = \mathcal{L}_H + \mathcal{L}_{DBI},$$

(13)
where

\[ \mathcal{L}_H = \frac{1}{24K} H_{\mu\nu}H^{\mu\nu}, \]  

(14)

with \( K = 8\pi G \) and the normalization fixed by string theory compactification. Since the primary issue is structure formation, which takes place before the onset of the accelerated expansion, we evaluate \( \mathcal{L}_{\text{DBI}} \) using the Newtonian gravity approximation and the light-cone gauge \([23, 24]\) for the D-brane embedding:

\[ X^i = x^i, \quad \frac{(X^0 + X^5)}{\sqrt{2}} = t, \quad \frac{(X^0 - X^5)}{\sqrt{2}} = \theta. \]

Then \( g_{00}^{(\text{ind})} = 2(\Phi + \dot{\theta}), \ g_{0i}^{(\text{ind})} = \theta_i, \ g_{ij}^{(\text{ind})} = -a^2\delta_{ij}, \) where \( \Phi \) is the peculiar gravitational potential, and these judicious choices yield

\[ \mathcal{L} = \frac{1}{4K} \left[ \frac{\dot{B}^k \dot{B}^k}{a^4} - \left( \frac{B^k}{a^6} \right)^2 \right] - \sqrt{A \left( 1 + \frac{B^k B^k}{a^4} \right)} \left( 2\dot{\theta} + 2\Phi + \theta_i \dot{\gamma}^{ij} \theta_j \right), \]

(15)

with

\[ \gamma^{ij} = \frac{a^2 \delta^{ij} + B^i B^j / a^2}{a^4 + B^k B^k} \]

(16)

being the symmetric part of the inverse of \( a^2 \delta_{ij} - \epsilon_{ijk} B^k. \)

The mass density is identified as \( \rho = -\partial \mathcal{L} / \partial \Phi \) which can be recast as a Bernoulli equation

\[ \dot{\theta} + \Phi + \frac{1}{2} \theta_i \gamma^{ij} \theta_j = \frac{A(1 + B^k B^k / a^4)}{2\rho^2}. \]

(17)

The peculiar gravitational potential is determined by Poisson’s equation

\[ \Phi_{,ii} = \frac{K}{2} a^2 (\rho - \bar{\rho}), \]

(18)

and the field equation for \( \theta \) yields the conservation equation

\[ \frac{1}{a^3} \frac{d}{dt} (a^3 \rho) + (\rho \gamma^{ij} \theta_j),_i = 0. \]

(19)

Finally, the field equation for \( B^i \) is

\[ \frac{1}{a^3} \frac{d}{dt} \left( \frac{\dot{B}^i}{a} \right) - \frac{B^j}{a^6} + \frac{2K \dot{A} B^i}{a^4 \rho} + 2K \rho \frac{\theta_i \theta_j}{a^2} - \gamma^{ij} \theta_j B^i / a^4 + B^j B^j = 0. \]

(20)

If we make the decomposition \( B^i = B^i_\perp + B^i_\parallel \) with the transverse part satisfying \( \partial_i B^i_\perp = 0, \) the key point becomes evident: whereas the longitudinal part \( B^i_\parallel \) suffers the same problem as Eq. (9), the transverse part does not experience spatial gradients.

Next, let us see the implications for linear perturbation theory. In addition to \( \delta = (\rho - \bar{\rho}) / \bar{\rho} \) and \( \psi = \theta - \bar{\theta}, \) we count \( B^i B^j / a^4 \) as first order because the Kalb-Ramond field, as a nonzero background, would be incompatible with isotropy. First, we show that once the perturbations enter the causal horizon, the longitudinal part of \( B^k B^k / a^4 \) in Eq. (17) undergoes damped oscillations, similar to those experienced by \( \delta \phi \) discussed in Sec. 2. Retaining the dominant terms, Eq. (20) becomes

\[ \frac{1}{a^3} \frac{d}{dt} \left( \frac{\dot{B}^i}{a} \right) - \frac{B^j}{a^6} + \frac{2K \dot{A} B^i}{a^4 \bar{\rho}} = 0. \]

(21)
As $H^2 = (\dot{a}/a)^2 = K\dot{\bar{\rho}}/3$, the last term in Eq. (21) is of order $H^2 A/\bar{\rho}^2$, so being negligible compared with the first term which is $O(H^2)$, until $a \sim 1$. Then, Eq. (21) simplifies to

$$\frac{d}{dt} \left( \frac{\dot{B}^i}{a} \right) - \frac{B^i_{ji}}{a^3} = 0, \quad (22)$$

In the transverse components of this equation the last term is absent, so the transverse solution reads

$$B^i_\perp(a, \vec{x}) \simeq c^i(\vec{x}) \int_0^a \frac{da}{H} = \frac{2}{5} \frac{c^i(\vec{x})}{H_0 \Omega^{1/2}} a^{5/2}, \quad (23)$$

where $\Omega$ is the equivalent matter fraction at high redshift and $c^i(\vec{x})$ are arbitrary functions of $\vec{x}$. Thus,

$$\frac{B^i_\perp B^i_\perp}{a^4} = \frac{4}{25} \frac{c^i c^i}{H_0^2 \Omega} a \quad (24)$$

grows linearly with the scale factor.

The longitudinal component of Eq. (22) in $k$-space reads

$$\frac{d}{dt} \left( \frac{\dot{B}_\parallel(k)}{a} \right) + \frac{k^2}{a^5} B_\parallel(k) = 0. \quad (25)$$

with the solution

$$B_\parallel(k) = a^{5/4} J_{5/2}(kd_c), \quad (26)$$

Hence $B_\parallel$ oscillates inside $d_c = 2\Omega^{-1/2} H_0^{-1} a^{1/2}$, with the amplitude increasing linearly with $a$. As a consequence, the longitudinal term $B^2_\parallel/a^4$ that enters the right-hand side of Eq. (17) oscillates with an amplitude decreasing as $a^{-2}$ compared with the transverse term $B^i_\perp B^i_\perp/a^4$ that grows linearly.

Since $a^3\dot{\bar{\rho}}$ is constant and $B^2_\parallel/a^4$ is damped inside $d_c$, retaining only the transverse part in (17), we obtain

$$\dot{\psi} + \Phi = \frac{A}{2\dot{\bar{\rho}}^2} \left( \frac{B^k_\parallel B^k_\perp}{a^4} - 2\delta \right), \quad (27)$$

$$\dot{\delta} + \frac{1}{a^2} \dot{\psi}_{,ii} = 0. \quad (28)$$

Owing to Eq. (24) we can arrange nonadiabatic perturbations for overdensities only such that $B^k_\parallel B^k_\perp/a^4 - 2\delta = 0$ with the assurance that they will hold independent of scale until $a \sim 1$. Combining Eqs. (18), (27), and (28) we find

$$\ddot{\delta} + 2H\dot{\delta} - \frac{3}{2} H^2 \delta = 0; \quad \delta > 0, \quad (29)$$

which is our main result: the growing mode overdensities here do not display the damped oscillations of the simple Chaplygin gas below $d_c$, but grow as dust. We remark that it matters little that this applies only for $\delta > 0$ since the Zel’dovich approximation implies that 92% ends up in overdense regions. Still, it should be kept in mind that underdensities will not behave as in [18].
4 Discussion

Clearly, there is an open question as to what inflation model can produce the initial conditions in the Kalb-Ramond field and brane embedding to allow subsequent structure formation. We believe this issue is likely to be closely related to another: namely, how does the Chaplygin-Kalb-Ramond model fit into the braneworld picture? We emphasize that $g^{(\text{ind})}_{\mu\nu}$ differs from $g_{\mu\nu}$ describes a brane different from that we inhabit, with the D3 Chaplygin brane moving rapidly until the onset of acceleration. Codimension-one braneworlds have little space for that.

The light-cone gauge formalism used here can be taken as a starting point for investigating questions beyond linear theory. In particular, one might hope that the acoustic horizon does resurrect at very small scales to provide the constant density cores seen in galaxies dominated by dark matter. Also, the “magnetic” nature of the Kalb-Ramond field, reflected in the cometric $\gamma_{ij}$, will generate rotation - it is pertinent to note that the simple, nonrotating, self-gravitating Chaplygin gas has a scaling solution $\rho(r) \propto r^{-2/3}$, far different from the $\rho(r) \propto r^{-2}$ wanted for flat rotation curves.

Ultimately, the model must be confronted with large-scale structure and the CMB. This necessitates a different gauge choice for $g^{(\text{ind})}_{\mu\nu}$. For the simple Chaplygin gas where the standard model feels the metric $g_{\mu\nu}$, it has been shown [25] that a good fit to the data is obtained if a vanishing sound speed is imposed on the Zel’dovich fraction. New quartessence brings in an added feature: the standard model can be considered as a non-Abelian Born-Infeld theory corresponding to a stack of coincident D3 branes, identifying the Abelian factor with $A_\mu$. Then the standard model experiences not the closed string metric $g_{\mu\nu}$, but rather the open string metric [26] $\gamma_{\mu\nu} = g_{\mu\nu} + B_{\mu\alpha}B_\alpha^\nu$ constructed as the inverse of the symmetric part of $(g + B)^{-1}$. This is reminiscent of the Einstein-Strauss nonsymmetric gravity theory, albeit the governing action is different and so the problems detailed in [27] may well be avoided.

Finally, the Kalb-Ramond field alone does not provide for transient acceleration as required by string theory. Rather than invoke an extra field as was done in [21], one may note that there is little difference between the Chaplygin gas D brane where $\sqrt{A}$ is constant, and tachyon models describing unstable D branes where $\sqrt{A} \to V(\theta)$ provided the potential is sufficiently flat [28] while still providing for eventual deceleration. The required flatness can be obtained through warped compactification\(^1\), which also explains the small value $A^{1/8} \sim \text{MeV}$.

5 Conclusions

In this paper we have shown how the Chaplygin gas model, when augmented by string theory features, can realize the quartessence scenario proposed earlier [4]. No ad-hoc constructions (such as generalized branes or generalized gravity) are required. To conclude, we would echo the sentiments of [18]: the day of quartessence cosmology is not at an end, but rather at a new beginning.

\(^1\)See, e.g., [29]. Note that the limits obtained there derive from the assumption that the photon field lives on the unstable brane, contrary to the picture here.
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