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Fluctuations in strongly coupled cosmologies & fundamental constant variations

Silvio A. Bonometto[1,2,3]

[1] Department of Physics, Astronomy Unit, Trieste University, Via Tiepolo 11, I 34143 Trieste, Italy, e-mail: bonometto@oats.inaf.it
[2] Istituto Nazionale di Astrofisica, Osservatorio Astronomico di Trieste, Via Tiepolo 11, I-34131 Trieste, Italy
[3] I.N.F.N., Sezione di Trieste, Via Valerio, 2 I 34127 Trieste, Italy

Abstract. In the early Universe, a dual component made of CDM and a scalar field Φ, exchanging energy, may naturally fall onto an attractor solution, making them a stationary fraction of cosmic energy during the radiation dominated era, as both such components expand as a^{-4} (as radiation), if their coupling $\beta > \sqrt{3}/2$. On this tracker solution they have early density parameters $\Omega_d = 1/4\beta^2$ and $\Omega_c = 2\,\Omega_d$ (field and CDM, respectively). In a previous paper it was shown that, at background level, this scenario naturally evolves towards a picture consistent with today's Universe, provided that a further component, expanding as a^{-3}, breaks the stationary expansion at $z \sim 3$–$5{\times}10^3$. Here we study the evolution of fluctuations of these background densities. Out of horizon fluctuation modes are determined. Their entry into horizon is also numerically evaluated as well as the dependence of Meszaros effect on the coupling intensity β. The transfer functions deduced in this approximation are quite consisten with data. This class of models then appears fully viable, with the clear advantage of predicting a 2–component DM, possibly easing the CDM crisis. A problem outlined here, however, concerns the possibility that this class of models induces a variation of fundamental constants, if the residual baryon–CDM coupling is too strong.

Key words. cosmology: theory, dark matter, dark energy, gravitation; methods: numerical

1. Introduction

Quite a few coincidences still lack an explanation, in cosmological models including DE (Dark Energy): (i) Why do we live the only era when DE and matter have similar weight? (ii) Why inhomogeneities are normalized so that DE allows fluctuations to form large scale non linear structures, then stopping any further density growth?

At variance from some belief, none of these paradoxes is eased if a self–interacting scalar field Φ replaces false vacuum, as suggested, e.g., by Ratra & Peebles (1988); Brax & Martin (1999, 2001); Brax et al. (2000); Baccigalupi et al. (2000); Bertolami & Martins (2000). In the former case, the self–interaction potential $V(\Phi)$ ought to contain a scale finely tuned by $\sim$ 30 o.o.m. in respect to the Planck scale m_p ($\sim 10^{28}$eV); tuning is somehow *worsened* in the latter case, when $\Omega_\Lambda^{1/4}$ ($\sim 10^{-2}$–10^{-3}eV) ought to be finely tuned

Send offprint requests to: S.Bonometto by $\sim$ 60 o.o.m.; however, in both cases, there are other physical scales similarly tuned: e.g., the EW transition scale and the neutrino mass difference scale.

However, it DE is a field Φ, it is natural to consider a coupling within the dark sector, as suggested by Ellis et al. (1989); Wetterich (1995); Amendola (1999); Amendola et al. (2002); Amendola et al. (2003) and debated by many authors; this may ease the (i) coincidence paradox, here above. It makes sense to replace false vacuum by a scalar field Φ, therefore, only if Φ couples to DM (Dark Matter).

Also in this case, however, it makes little sense to debate about the shape of $V(\Phi)$. As a matter of fact, observations scarcely constrain the very DE state equation $w(a)$ (a : scale factor). The very Euclid[1] (Laureijs et al., 2011) experiment is expected to constrain $w(a)$ derivative at $z = 0$ with an error $\sim$ 20 %, as shown by Joachimi & Briddle (2010). Constraining $V(\Phi)$ is even harder and parameters inside it can only be constrained after we assume a specific shape; this was done, e.g., by La Vacca & Kristensen (2009).

Accordingly, here we shall debate cosmologies where DE to be a scalar field interacting with DM, but will try to bypass the choice of a potential, by focusing on parameters closer to observations.

The point is that DM–DE interactions allow DE to be fed fresh energy, so keeping a significant cosmic component at any redshift, not only in our era, as shown, e.g., by Amendola (1999). A fairly large DM–DE interaction scale

$$C = b/m_p = \sqrt{16\pi/3}\ \beta/m_p \tag{1}$$

is consistent with data. In fact b values up to $\sim$ 0.4 are allowed, namely if neutrinos have a non–negligible mass; this point, first outlined by La Vacca et al. (2009), was then confirmed, from other points of view or using different datasets by Xia (2009); Kristiansen et al. (2010); Pettorino et al. (2011). In this way, DE is allowed to keep at the $\sim$ 1 % level in respect to CDM up to our era, whose eve is however still characterized by a specific event: the transition of the DE field from the kinetic to the potential regime.

[1] http://www.euclid-ec.org

In fact, pressure and energy density of a scalar field read

$$p_d = \dot\Phi^2/2a^2 - V(\Phi)\ , \quad \rho_d = \dot\Phi^2/2a^2 - V(\Phi) \tag{2}$$

(differentiation is in respect to conformal time) so that the state parameter $w = p/\rho$ is $\sim +1$ [~ -1] when the $\dot\Phi^2/2a^2$ [$V(\Phi)$] term dominates. In the former case, $\rho_d \propto a^{-6}$ would rapidly become negligible unless CDM transfers energy on DE, so diluting at a rate (slightly) faster that a^{-3}. Agreement with data has been found by keeping this decreased rate as close as possible to a^{-3}, and this is why DE is however significant, but cannot exceed $\sim$ 1 % of CDM density.

In order that DE becomes significant, as it is today, the progressive Φ increase ought to cause a (recent) transition to potential dominance. The energy transfer from CDM keeps quite low, but ρ_d relevance rapidly increases because of the progressive fall of w.

2. Strongly coupled DE

Background properties

An alternative to this picture has been recently considered in the literature (Bonometto, Sassi & La Vacca 2012). If the coupling strength is drastically increased, both the CDM density ρ_c and the (kinetic) DE density ρ_d can dilute $\propto a^{-4}$, as radiation does. What Bonometto, Sassi & La Vacca (2012) find is that this regime is an attractor, provided that $\beta^2 > 3/4$, while CDM and DE have then early density parameters

$$\Omega_c = 1/2\beta^2\ , \quad \Omega_d = 1/4\beta^2\ . \tag{3}$$

This regime, characterized by a stationary density sharing, approaches the observational picture if another DM component breaks it, by overcoming the density of the “radiative” components.

At variance from Bonometto, Sassi & La Vacca (2012), here we consider the option that this latter uncoupled DM component is warm. Figure 1 then describes the evolution of densities for $\beta = 3.5$ and 7 . In the plots, the redshift $z_\pm$ (DE turning from kinetic to potential) is suitably tuned to yield a density parameter

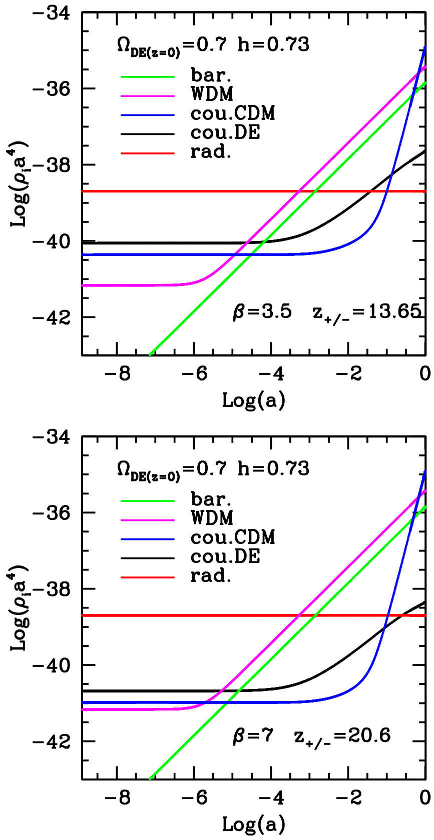

Fig. 1. Evolution of background components in cosmologies with coupled CDM and uncoupled WDM (thermal particles of mass $m_w = 0.1$ keV, 2 spin states, effective temperature $0.25\,T_\gamma$). The cases $\beta = 3.5$ and 7 are considered, both yielding $\Omega_d = 0.7$ when the redshift $z_\pm$ (DE turning from kinetic to potential) is suitably tuned.

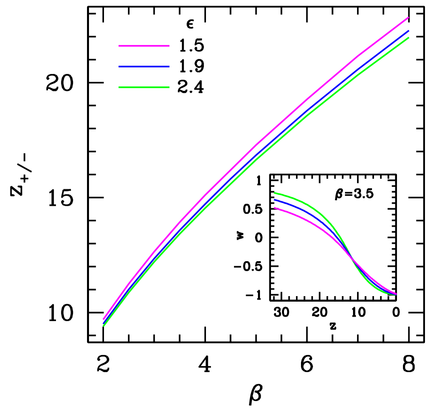

Fig. 2. Transition from $w = +1$ to $w = -1$ of DE state equation; $h = 0.73$ and $\Omega_d = 0.7$ at $z = 0$ are required. The redshift when $w = 0$ ($w_\pm$) exhibits just a mild dependence on the transition laws, sampled by the ϵ parameter and shown in the inner box. WDM as in previous Figures.

$\Omega_d = 0.7$ for DE, at $z = 0$. Results are very mildly dependent on the shape of the kinetic–potential transition, as shown in Figure 2. In its inner frame we show the shapes of the transitions, parametrized by ϵ.

This confirms that the properties of the DE self–interaction potential are hard to be observationally constrained.

Fluctuations

Let us now consider density fluctuations on this background (the detailed computations leading to the results herebelow will be presented elsewhere). Before the entry into the horizon, in a synchronous gauge where the metric reads

$$ds^2 = a^2(\tau)\left[d\tau^2 - (\delta_{ij} + h_{ij})dx^i dx^j\right] , \qquad (4)$$

metric perturbation can then be expanded as follows:

$$h_{ij}(\tau, \mathbf{x}) = \int d^3k\, e^{i\mathbf{k}\cdot\mathbf{x}}[n_i n_j h(\tau, \mathbf{k}) + \\ + (n_i n_j - \delta_{ij}/3)\, 6\eta(\tau, \mathbf{k})] \qquad (5)$$

with $\mathbf{k} = \mathbf{n}k$. Einstein equations then yield

$$\ddot{h} + (\dot{a}/a)\dot{h} = -(8\pi/m_p^2)\,(\delta\rho + 3\delta p)\, . \qquad (6)$$

The gravity sources are: (i) radiation, for which $\delta\rho + 3\delta p = 2\rho_r\delta_r$; (ii) baryons; (iii) uncoupled CDM or WDM; (iv) coupled CDM; and, finally, (v) the DE field Φ, for which

$$\delta(\rho_\phi + 3p_\phi) = \delta\left[4(\Phi_1^2/2a^2) - 2V(\Phi)\right] = \\ = \bar{\Phi}_1\,\dot{\phi}/a^2 - 2V'(\bar{\Phi})\,\phi \qquad (7)$$

being $\Phi = \bar{\Phi} + \phi$ (here $\bar{\Phi}$ is the background field; however, in the sequel, we shall continue to omit the bar on it).

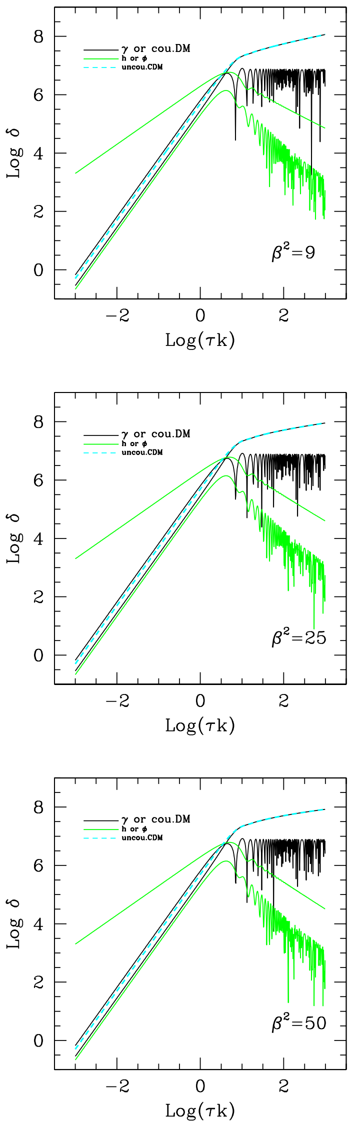

Fig. 3. We compare early fluctuation evolutions in 4 cosmologies. 3 of them are shown in this Figure and are strongly coupled with $\beta \simeq 3,\ 5$ and ~ 7.1 (from top left to bottom left). Uncoupled and coupled DM fluctuations overlap as soon as their wavelength is reached by horizon. Meszaros effect exhibits a slight dependence on β, which will translate into (slightly) different effective slopes in the transfered spectra. For the sake of comparison, in the next Figure we show the case of a standard LCDM cosmology.

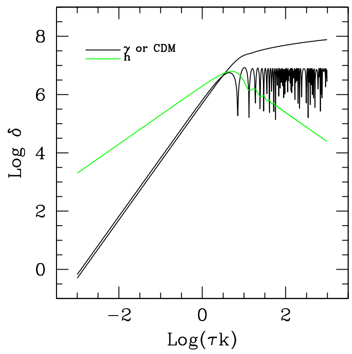

Fig. 4. See previous Figure.

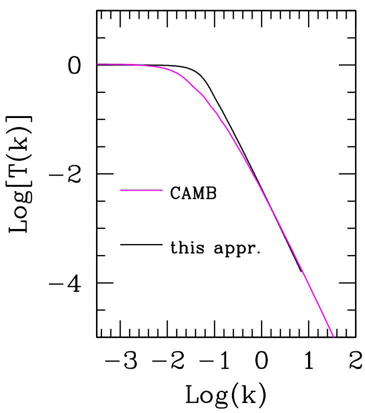

Fig. 5. Transfer function for a ΛCDM model, obtained according to our simplified algorithm, compared with the transfer function obtained from CAMB, for the same model. In the lower frame fluctuation evolutions are shown.

By assuming then, out of the horizon,

$$h = A\tau^a\ ,\quad \delta_r = R\tau^r\ ,\quad \delta_w = W\tau^w\ ,\quad \delta_c = M\tau^c\ ,$$

$$\phi = \frac{m_p}{b}\varphi \ \text{ and } \ \varphi = F\tau^f\ , \tag{8}$$

we however find $r = w = c = a \equiv x$ and two sets of modes:
modes (a): $x = \pm 2$, being β independent;
modes (b): $x = (1/2)[-1 \pm 3^{1/2}(1/\beta^2 - 1)^{1/2}]$.

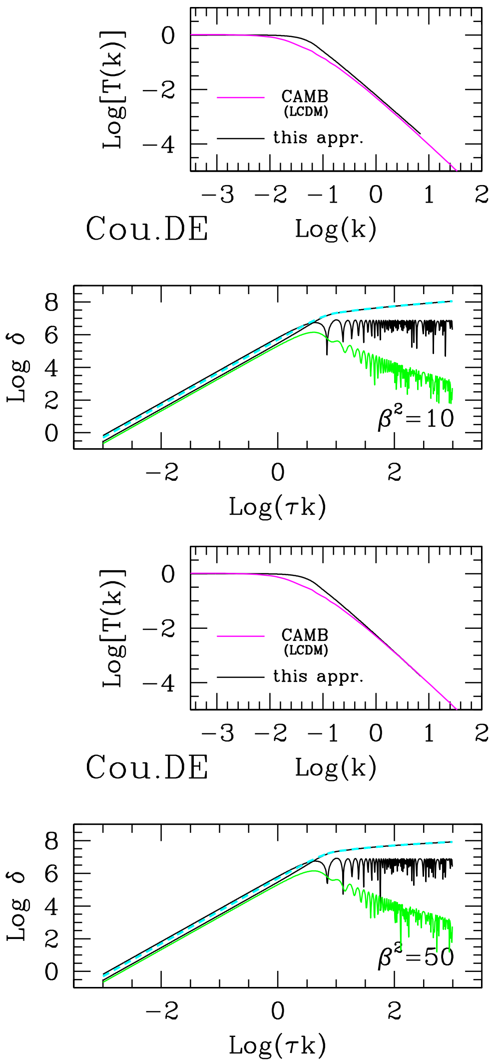

Fig. 6. As previous plot, for two coupled DE cosmologies. A slight dependence of the transfered spectrum slope on β is easily predicable, with greater β's better approaching ΛCDM.

For all modes we find $R = -(2/3)A$ and $W = (-1/2)A$, also implying $R = (4/3)W$, as expected. In the case of the increasing (a) mode, we have

$$M = -(3/14)A\ , \quad F = -(2/7)A\ , \qquad (9)$$

while $M = (3/7)W$. Fluctuations, therefore, are less than half of uncoupled–DM. In coupled DE models coupled–CDM fluctuations are mostly enhanced, in respect to uncoupled components, by a β dependent factor. In initial conditions an opposite (β–independent) behavior is met. Modes (b) grow quite slowly and will be discussed elsewhere.

With initial conditions set in accordance with an increasing (a) mode, we find the fluctuation evolution in Figure 4. Here 3 values for β are considered, and the program used to study fluctuation evolution is also applied to a standard ΛCDM case, for the sake of comparison.

The program assumes a tight baryon–photon binding yielding a fluid with speed of sound $c_s \equiv 1/\sqrt{3}$, no neutrinos, no recombination, and $\dot a/a \equiv 1/\tau$. What is clearly detected, however, is Meszaros effect, the basic mechanism yielding the shape of the transfer function. By using it, we can provide an estimate of the transfer function itself.

In Figure 5, the transfer function found in this way is compared with CAMB transfer function (normalized to unity for $k \to 0$), in the case of no baryons. For coupled DE, in the same way, we obtain the transfer functions shown in Figure 6 and compared there with the previous CAMB result.

Our simplified numerical approach does not allow more detailed predictions, but allows us to conclude that, for a suitable cosmological parameter choice, strongly coupled–DE cosmologies are at least as efficient as ΛCDM to fit observational data.

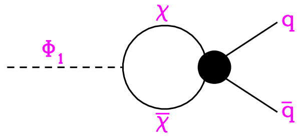

Fig. 7. Coupled–DE quark interaction diagram ($\Phi_1 = \dot\Phi$). This diagram could interfere with the quark–Higgs coupling setting quark and lepton masses.

3. Fundamental constant variations?

Rather, a problem for these cosmologies is schematized in Figure 7. The lagrangian coupling needed to allow energy transfer from DM to DE, shown in Bonometto, Sassi & La Vacca (2012) (see also Das, Corasaniti & Khoury, 2006), is likely to imply an indirect coupling of quarks (q in the Figure) with the DE field. In the figure χ are coupled–DM particles and the $\dot{\Phi}$–χ vertex yields a β factor in the rate of the process. The rate is also weighted by the unknown χ–q interaction, which might well exist, even though the m.f.p. of χ's, in the present epoch, exeeds the horizon by lots of orders of magnitude.

Notice that $\dot{\Phi} \propto (1+z)$, so that the diagram might acquire a greater importance in the past. This diagram could interfere with the quark–Higgs interactions, setting quark and lepton masses. In a forthcoming work we shall provide the limits on the χ–q effective interaction, for gradually increasing β values, required in order that fundamental constants are substantially z independent.

4. Conclusions

Let us finally remind a few inconsistencies of the ΛCDM model, on sub-galactic scales, put in evidence by N–body simulations, namely if DM is assumed to be "cold". A first difficulty concerns the amount of substructure in Milky Way sized haloes, as shown by Klypin et al. (1999) and Moore et al. (1999). Models involving CDM overpredict their abundance by approximately one order of magnitude. A second issue concerns the density profiles of CDM haloes in simulations, exhibiting the typical NFW cuspy behavior, as confirmed by Moore et al. (1994); Flores & Primack (1994); Diemand et al. (2005); Macciò et al. (2007) and a few other authors; on the contrary, the density profiles inferred from rotation curves suggest a core like structure (de Blok et al., 2001; Kuzio de Naray et al. 2009; Oh et al. 2011). A third issue concerns dwarf galaxies in large voids: Tikhonov et al. (2009), Zavala et al. (2009), and Peebles & Nusser (2010) recently re–discussed their abundance; a substantial excess is likely to have been observed, although no full agreement has yet been achieved.

A streaming length of the DM component consistent with a thermal relic of particles with mass $\sim$ 2–3 keV seems to yield better predictions. There is a number of candidates for such WDM, as sterile neutrinos or gravitinos. WDM leads to a suppression of the linear power spectrum on galactic and sub–galactic scales, as first seen by Bonometto & Valdarnini (1984) and solves several above problems, as debated by many authors, e.g. Hogan & Dalcanton (2000); Viel et al. (2005); Abazajian & Koushiappas (2006). In particular, the profiles of WDM haloes, similar to CDM haloes in the outer regions, flatten in the inner regions, as predicted by Villaescusa-Navarro & Dalal (2011) and found in simulations (Colin et al. 2008; Macciò et al. 2012).

However, the core size for thermal candidates allowed by large scale constraints (Lyman–α and lensing), is 30–50 pc, while observed dwarf galaxies exhibit cores of $\sim$ 1000 pc's (Walker & Penarrubia 2011). Such a dwarf galaxy core would require masses $<$ 0.1–0.3 keV; but such a warm candidate yields a streaming length exceeding the size of these very dwarf galaxies (Macciò & Fontanot 2010; Macciò et al. 2012).

In view of these difficulties, some authors have proposed that a large amount of WDM is accompanied by a smaller amount of CDM. The WDM particle velocities could then be greater, while a low–mass population is however produced by CDM clustering. It is worth outlining that these suggestions have been put forward quite indipendently of any particle or cosmic model. In particular, assuming *ad hoc* a twofold dark matter component appears as a rather extreme supposition, certainly not easing coincidence problems. Such an assumption is made just because simpler models are apparently facing a deadlock.

These predictions cannot be soon translated to the present case, let alone because the CDM component keeps coupled to DE. All that does not prevent us from outlining a strict similarity between data requirements and this class of models.

In Figures 4 no free–streaming was considered. If WDM free streams, the shown dependence on $k\tau$ would no longer be k independent. What is however sure is that, even if large k values are considered, so that free streaming fully destroys primeval WDM fluctuations, they will be re–generated by coupled CDM fluctuations, after WDM derelativisation. This will occur at a redshift between 10^3 and 10^4, when CDM density is a fraction of WDM; the effect is similar to baryons falling onto DM seeds after recombination. The amplitude of the WDM–baryon spectrum can then be expected to be smaller below the WDM streaming length, in a β dependent way and this effect could be used to estimate β.

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