Strongly Coupled Cosmologies

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Abstract

Models including an energy transfer from CDM to DE are widely considered in the literature, namely to allow DE a significant high–z density. Strongly Coupled cosmologies assume a much larger coupling between DE and CDM, together with the presence of an uncoupled warm DM component, as the role of CDM is mostly restricted to radiative eras. This allows us to preserve small scale fluctuations even if the warm particle, possibly a sterile ν, is quite light, $O(100\text{eV})$. Linear theory and numerical simulations show that these cosmologies agree with ΛCDM on supergalactic scales; e.g., CMB spectra are substantially identical. Simultaneously, simulations show that they significantly ease problems related to the properties of MW satellites and cores in dwarfs. SC cosmologies also open new perspectives on early black hole formation, and possibly lead towards unificating DE and inflationary scalar fields.

1. Introduction

We discuss a new family of models, starting from their features in the radiative eras. In such epochs, they are characterized by two extra components, in top of the usual radiative ones: a scalar field Φ and a peculiar CDM component, with energy densities and pressures $\rho_\Phi$ & $p_\Phi$ and $\rho_c$ & $p_c$, respectively. As we assume $\rho_\Phi \approx p_\Phi \approx \Phi^2/2a^2$ and being $p_c \approx 0$, it should be $\rho_\Phi \propto a^{-6}$, $\rho_c \propto a^{-3}$. It is not so, because the models assume an energy flow from CDM to the field, due to a Yukawa–like interaction Lagrangian

$$\mathcal{L}_I = -\mu f(\Phi/m)\bar{\psi}\psi ,$$

being an attractor for the system made by the equations of motions of Φ and ψ (m_p: the Planck mass, $\tau_\gamma$: a generic reference conformal time). Eqs. (3) and (4) imply that $\rho_c \propto \rho_\Phi \propto a^{-4}$, so that CDM and the field dilute at the same rate of the radiative components. On the attractor, the constant early state parameters $\Omega_c$ and $\Omega_\Phi$ (CDM and Φ, respectively) shall then read

$$\Omega_c = 1/(2\beta^2) , \quad \Omega_\Phi = 1/(4\beta^2)$$

with

$$\beta^2 = (3/16\pi)b^2 .$$

What happens is that the flow of energy from CDM to the field fastens (slows down) the dilution of CDM ($\Phi$); accordingly, the scale factor exponents change: for the former component from -3 to -4, for the latter one from -6 to -4. The notable point is that this behavior occurs along an attractor: if starting from generic initial conditions, with $\Omega_\Phi$ and/or $\Omega_c$ different from eq. (5), the e.o.m rebuild the conditions (5), also suitably synchronizing the rates of energy transfer and cosmic expansion. For a detailed proof, see Paper A, wherefrom Figure[1] is taken.
Figure 1: Starting from generic initial conditions, the attractor is soon recovered: both $\rho_c$ and $\Phi$ were initially “wrong” or, equivalently, the rate of expansion was “wrong”. Here $\beta = 2.8$, so that the overall density of CDM and $\Phi$ is smaller than half extra $\nu$ species.

Figure 2: Background evolution. WDM derelativisation breaks the former conformal invariance, allowed by CDM–$\Phi$ coupling. Different curves yield models where coupling either persists down to $z = 0$ or fades earlier (see text). The $w$ (field state parameter) shift, from $+1$ to $-1$ is tuned to account for a present DE density parameter $\Omega_d = 0.7$. These models are then characterized by three phases: (B) before (the radiative eras); (D) during (them); (A) after (matter–radiation equality).

The stages (D) and (A) were treated both in [1] and in a further Paper B [2], focused on fluctuation evolution, finding that, from the above attractor, the models naturally evolve towards a picture consistent with today’s Universe. Besides of baryons this requires a WDM (Warm Dark Matter) component (see, e.g., [3]).

More in detail: In the present epoch DM is essentially warm, $\Phi$ has turned into quintessential DE (Dark Energy), while CDM, keeping an almost negligible density, seems to play just an ancillary role.

In Figure 3 we show the background evolution in some spatially flat models with $\Omega_d = 0.7$, $\Omega_b = 0.045$, $h = 0.685$, in agreement with Planck results. Model dynamics is somehow reminiscent of the coupled DE op-

tion, studied by many authors [4]. Here, however, coupling plays its key role through radiative eras. Switching it off (or letting $m \to \infty$) after WDM has derelativized, could even ease the fit with data. Besides of a coupling persistent down to $z = 0$, we therefore consider the cases of $\beta$ fading exponentially at $z = z_{\text{der}} \times 10^{-D}$ (here we assume that WDM is a sterile $\nu$ with a former thermal distribution and the redshift $z_{\text{der}}$ is when $m_\nu = T_\nu$; the exponent $D$ is dubbed delay).

2. Fluctuation evolution

In Paper B it is widely discussed how the public program CMBFAST (or, similarly, CAMB) is to be modified to follow fluctuation evolution in these models. Changes involve also out–of–horizon initial conditions.

A notable feature of these models is that there is no cut in the transfer function of WDM. The point is that, in the non–relativistic regime, the presence of a CDM–$\Phi$ coupling yields an increase of the effective self–gravity of CDM, as though

$$G \to G' = G(1 + 4\beta^2/3)$$  \hspace{1cm} (7)

(see [5]). The interaction of CDM with other components is set by $G'$, the $G'$–shift concerning just CDM–CDM gravity. When fluctuations reach the horizon, the CDM density parameter is $O(\beta^2)$. However, as its self interaction is boosted by a factor $O(\beta^2)$, it evolves as though $\Omega_c \sim 2/3$, independently of $\beta$.

The other, velocity dependent, changes in CDM dynamics, not discussed here, do not modify the fact that the growth of CDM fluctuations is never dominated by the baryon–radiation plasma. Accordingly, while
baryons and radiation yield sound waves, and collisionless components suffer free streaming, the CDM fluctuations $\delta_\ell$ suffer no Meszaros' effect and steadily grow. Later on, when WDM derelativizes and/or baryons decouple from radiation, $\delta_\ell$ has grown so large, to cause the re–generation of fluctuations in the other components, and this is an essential feature to meet specific data ΛCDM models do not fit.

In Figure 4 we give an example of the linear evolution of density fluctuations down to $z = 0$. The spectrum

$$\Delta^2(k) = \frac{1}{2\pi^2} k^3 P(k) \quad \text{with} \quad P(k) = \langle |\delta|^2 \rangle$$

(8)
is then shown in Figure 5 at $z = 0$, as obtainable from the linear code, and with arbitrary normalization. The model considered has $D = 2$; CDM–CDM interactions are therefore ruled by the normal $G$ since a redshift $z \sim 10^3$. Let us outline, first of all, that different components exhibit different spectra, even at $z = 0$, with CDM showing the most peculiar behavior. While, in the relativistic regime, outside the horizon, CDM fluctuations behave similarly to other components, later on, between the entry in the horizon and $z \simeq 10^{-2} z_{der}$, they undergo an authonomous fast growth. The entry in the horizon and, therefore, the duration of this growth depends on mass scale. Figure 4 shows a CDM spectral function exceeding the total function by one o.o.m. already at $k \simeq 40 \, h{\text{Mpc}^{-1}}$, corresponding to a mass scale $M \simeq 9 \times 10^9 M_{\odot} h^{-1}$.

The distinction between baryon and CDM spectra, in Figure 5 however, is artificial, being due to the algorithm used. After CDM–Φ decoupling baryons and CDM fulfill the same linear equations, so that possible non–linearities depend on the amplitude of the weighted sum $\Omega_\delta \delta + \Omega_\phi \delta_\phi$. Of course, there are mass scales where CDM is already non–linear at $10^{−D_{der}}$, for any $D > 0$; for the model in Fig. 4 this occurs for $M < 10^{6} M_{\odot} h^{-1}$. If CDM non–linearities arise, the code is scarcely predictive, even though we may conjecture its results to be reasonably well approximated, for the non–CDM components, until $\delta_\ell < 10$ (approximately a spherical fluctuation turnover). This does not necessarily mean that the model is unphysical.

3. Dwarf galaxy cores and MW satellites

The key issue, however, is the similarity between this model and ΛCDM. Plotting $\Delta^2(k)$ aims to stress model discrepancies, less evident in transfer functions or integral quantities (e.g. $\sigma_8$). ΛCDM is a benchmark for any model improvement, its main discrepancies from data being: (i) the number of MW satellites [6], (ii) (dwarf) galaxy cores [7]. To overcome such difficulties, the option of replacing cold with warm DM has been explored. According to [8], however, dwarf galaxies in ΛWDM N–body simulations exhibit a core radius

$$R_{\text{core}} \sim 1 \, h^{-1} \text{kpc} \left(100 \, \text{eV} / m_\nu\right)^{1.8}.$$  

(9)

Observations require $R_{\text{core}} \sim 0.5–1 \, h^{-1} \text{kpc}$; ΛWDM cosmologies, therefore, are no solution unless $m_\nu \sim 100$ eV. In Figure 5 we then show linear spectral functions for a number of $m_\nu$ values, showing that a WDM cosmology with $m_\nu \sim 100$ eV yields no systems even with mass $\sim 10^{13} M_{\odot} h^{-1}$. In the literature, it is often assumed that cosmological data suggest a sterile $\nu$ with mass $m_\nu \sim 2–3 \, \text{keV}$. According to Fig. 5 such $\nu$ is barely sufficient to produce structures over galactic scales, but yields no improvement on the above problems (i),(ii). Within SC cosmologies, on the contrary, we predict a sterile neutrino with mass $\mathcal{O}(100 \, \text{eV})$. 

Figure 4: Linear spectral function in a model with $D = 2$. The scale in ordinate is arbitrary. The dashed thin line is the spectral function for a ΛCDM model with the same normalization.

Figure 5: Linear spectral function in ΛWDM models with various neutrino masses, as indicated in the frame.
Figure 6: Profile of a bound system of mass $\sim 6 \times 10^{10} M_{\odot} h^{-1}$ in N–body simulation of a SC cosmology (in blue; here $m_p = 90$ MeV) compared with the same system in a $\Lambda$CDM model (in black).

N–body simulations for these cosmologies are in progress and suggest that, besides of yielding a fair core profile (see preliminary results in Figure 6) such models approximately half the prediction of MW satellites.

4. Discussion

Besides of easing these problems, SC cosmologies open two significant perspectives. Let us then briefly recall that a cosmological constant $\Lambda$ is due to the (slow) CDM fluctuation growing even in the presence of Meszaros’ effect. SC cosmologies are surely characterized by a much more rapid CDM fluctuation growth. In any SC model, therefore, CDM fluctuations on scales $< \sim 10^4 M_{\odot} h^{-1}$ get non linear before radiative expansion ends. The rate of a spherical collapse is not reached, it could approach a relativistic regime.

This suggests us two specific comments: (i) during pre–relativistic non–linear stages, the coherence between CDM and other component distributions might fade over small scales, causing a low mass cut–off in the transfer function, however much below the one in Figure 5; (ii) when approaching a relativistic regime, the approximated expression (7) looses validity. Let us recall that a cosmological constant $\Lambda$ was already found to increase the threshold for cosmological BH formation (9). The problem here is more intricate, however: a $\Lambda$ component is already significant at high $z$; a possible conjecture is that BH formation is also suppressed, at least until $\beta$–coupling is active. The time and space dependence of $\Lambda$, however, could source unexpected effects and, although equations have been set (9), only their numerical treatment can provide a reliable answer. According to the conjecture, however, the fading of $\beta$–coupling could then trigger BH production.

Let us finally comment about the stage B (before the onset of radiative eras). SC cosmologies however set a bridge between the late inflationary stages and our epoch. During this period, the $\Phi$ field underwent just a logarithmic growth, its present value being just $O(60)$ times its value at inflation end. Meanwhile, kinetic energy has decreased by $\sim 240$ o.o.m.’s. The transition of DE state parameter $w$ from 1 to -1, therefore, can be due to the same potential causing the inflationary expansion only if it exhibits an exponential dependence on $\Phi$. This is somehow reminiscent of the interaction $L_I$ (eq. 1) which, anyhow, cannot be responsible for the potential energy. The shape of $L_I$ however tells us that a large field inflation might occur when $\Phi \sim m_p \gg m$ (requiring then $b \gg 1$), so that the $\Phi$–$\psi$ interaction is practically switched off. A progressive decrease of $\Phi$ could then reactivate $L_I$ causing a reheating, when $\Phi$ turns into $\psi$ quanta, to finally stabilize on the attractor solution. It might be appealing to consider this perspective also within the frame of primeval particle production and annihilation due to curvature variation (10).

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