WIMPs during Reheating

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Abstract. Weakly Interacting Massive Particles (WIMPs) are among the best-motivated dark matter candidates. In the standard scenario where the freeze-out occurs well after the end of inflationary reheating, they are in tension with the severe experimental constraints. Here, we investigate the thermal freeze-out of WIMPs occurring *during* reheating, while the inflaton ϕ coherently oscillates in a generic potential $\propto \phi^n$. Depending on the value of n and the spin of the inflaton decaying products, the evolution of the radiation and inflaton energy densities can show distinct features, therefore, having a considerable impact on the freeze-out behavior of WIMPs. As a result of the injection of entropy during reheating, the parameter space compatible with the observed DM relic abundance is enlarged. In particular, the WIMP thermally averaged annihilation cross-section can be several magnitudes lower than that in the standard case. Finally, we discuss the current bounds from dark matter indirect detection experiments, and explore future challenges and opportunities.

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1 Introduction

The existence of non-baryonic dark matter (DM) in the Universe is very compelling, as suggested by both astrophysical and cosmological observations [1]. A viable DM candidate has to satisfy several properties: It has to be electromagnetically neutral, stable at the cosmological scales, and non-relativistic at the matter-radiation equality in order to allow structure formation. Finally, it must feature a relic density $\Omega_{\rm dm}h^2 \simeq 0.12$, accounting for 27% of the total energy budget of the Universe [2, 3].

The most prominent DM production mechanism in the early universe corresponds to the weakly interacting massive particle (WIMP) paradigm [4, 5]. In the standard WIMP scenario, DM has a mass at the electroweak scale and couples to the standard model (SM) thermal plasma with a sizable strength, typical of electroweak interactions. WIMPs reach thermal equilibrium with the SM thermal plasma and eventually *freeze out*, giving rise to the observed DM relic abundance. This DM freeze-out is commonly assumed to occur well after the end of reheating, when the Universe energy density is dominated by SM radiation. To match observations, a thermally averaged annihilation cross-section $\langle \sigma v \rangle = \text{few} \times 10^{-26} \text{ cm}^3/\text{s}$ is typically required [6]. The WIMP mechanism is particularly interesting because it could be tested in a number of complementary ways, including direct, indirect, and collider probes. However, current null experimental results and severe constraints on the natural parameter space motivate some quests beyond the standard WIMP paradigm [4].

There are several possibilities to evade the strong experimental constraints. For example, DM might have feeble interactions with the SM thermal plasma so that it is produced out of equilibrium by the so-called freeze-in mechanism (FIMP) [7–13]. Alternatively, WIMPs might have decoupled during a non-standard cosmological epoch [14], with an annihilation cross section much smaller than the one enforced in the standard case [15]. In general, the production of DM in scenarios with a non-standard expansion phase has recently gained increasing interest; see, e.g., Refs. [15–52]. For earlier work, see also Refs. [53–64]. Moreover, scenarios where freeze-out occurs during reheating are also viable, and in such cases, the freeze-out temperature $T_{\rm fo}$ is higher than the reheating temperature $T_{\rm rh}$; note that the latter characterizes the onset of the radiation-dominated era. We note that WIMP scenarios with low reheating temperature have been widely investigated in the literature [28, 60, 65–68], usually triggered by the decay of a long-lived scalar field with an equation-of-state (EoS) parameter $\omega = 0$ and

a constant decay rate Γ_{ϕ} .¹ Finally, we note that this kind of scenario in which DM reaches chemical equilibrium with the SM bath eliminates the possible overproduction of DM during inflation [83, 84].

In this work, different from the aforementioned literature, we investigate the phenomenology of WIMP freezing out during reheating, where the inflaton field ϕ oscillates around the minimum of a generic potential $\propto \phi^n$. Such potentials naturally arise in a number of inflationary scenarios like the α -attractor *T*-model [85, 86], or the Starobinsky model [87–90]. In such a case, both the inflaton EoS and its decay rate are modified compared to the usually assumed case (i.e. n = 2). In particular, the EoS depends on the shape of the potential $\omega = (n-2)/(n+2)$ [91], and Γ_{ϕ} develops a time dependency [92, 93]. As a consequence, the inflaton and SM energy densities exhibit a non-standard evolution during reheating [94], which has a strong impact on WIMP decoupling.² Depending on the value of *n*, the inflaton decay products, as well as the reheating temperature, we find that due to the large injection of entropy during reheating, $\langle \sigma v \rangle$ can be several orders of magnitude smaller than in the standard case, thus evading the experimental constraints. We highlight that some regions of the new parameter space are currently being tested by several experiments. Furthermore, we notice that the future CTA experiment might increase our ability to probe different scenarios.

The remainder of the paper is organized as follows. In Sect. 2 we revisit inflationary reheating, in particular, the evolution of the inflaton and the radiation energy densities, with an inflaton oscillating in a potential $\propto \phi^n$. In Sect. 3, we investigate WIMP DM production focusing on the case where freeze-out occurs *during* reheating. Our findings are summarized in Sect. 4.

2 Cosmology during Reheating

We consider an inflaton ϕ that, after the end of inflation, oscillates at the bottom of a monomial potential V of the form

$$V(\phi) = \lambda \frac{\phi^n}{\Lambda^{n-4}}, \qquad (2.1)$$

where λ is a dimensionless coupling and Λ an energy scale. Potentials with these types of minima naturally arise in a number of inflationary scenarios, such as the α -attractor Tmodel [85, 86], or the Starobinsky model [87–90]. The equation of motion for the oscillating inflaton field is given by [91]

$$\ddot{\phi} + (3H + \Gamma_{\phi})\dot{\phi} + V'(\phi) = 0,$$
(2.2)

where H denotes the Hubble expansion rate and Γ_{ϕ} the inflaton decay rate. The dots and primes correspond to derivatives with respect to time t and ϕ , respectively. By multiplying the above expression by $\dot{\phi}$, Eq. (2.2) can be rewritten as

$$\frac{d}{dt}\left[\frac{1}{2}\dot{\phi}^2 + V(\phi)\right] = -(3H + \Gamma_{\phi})\dot{\phi}^2.$$
(2.3)

¹For studies on baryogenesis with a low reheating temperature or during an early matter-dominated phase, see Refs. [65, 69–72] and [73–75], respectively. Furthermore, primordial gravitational wave production in scenarios with an early matter era has recently received particular attention [76–82].

²We note that the phenomenology of FIMPs [92, 93, 95-97] and the QCD axion DM [98] with a time-dependent decay rate has recently been analyzed.

By defining the energy density and pressure of ϕ as $\rho_{\phi} \equiv \frac{1}{2}\dot{\phi}^2 + V(\phi)$ and $p_{\phi} \equiv \frac{1}{2}\dot{\phi}^2 - V(\phi)$, together with the equation-of-state parameter $\omega \equiv p_{\phi}/\rho_{\phi} = (n-2)/(n+2)$ [91], Eq. (2.3) gives the evolution of ρ_{ϕ} as

$$\frac{d\rho_{\phi}}{dt} + \frac{6n}{2+n} H \rho_{\phi} = -\frac{2n}{2+n} \Gamma_{\phi} \rho_{\phi} \,. \tag{2.4}$$

During reheating, that is, in the range $a_I \ll a \ll a_{\rm rh}$, where a is the scale factor, and a_I and $a_{\rm rh}$ correspond to the scale factors at the beginning and end of the reheating, respectively, Eq. (2.4) admits the analytical solution

$$\rho_{\phi}(a) \simeq \rho_{\phi}(a_{\rm rh}) \left(\frac{a_{\rm rh}}{a}\right)^{\frac{6\,n}{2+n}}.$$
(2.5)

As during reheating the Hubble expansion rate is dominated by the inflaton energy density, it follows that

$$H(a) \simeq \begin{cases} H_{\rm rh} \left(\frac{a_{\rm rh}}{a}\right)^{\frac{3n}{n+2}} & \text{for } a \le a_{\rm rh} ,\\ H_{\rm rh} \left(\frac{a_{\rm rh}}{a}\right)^2 & \text{for } a_{\rm rh} \le a . \end{cases}$$
(2.6)

At the end of reheating (i.e., at $a = a_{\rm rh}$), the inflaton and radiation energy densities are equal, $\rho_R(a_{\rm rh}) = \rho_\phi(a_{\rm rh}) = 3 M_P^2 H_{\rm rh}^2$. Note that in order not to spoil BBN, the reheating temperature must satisfy $T_{\rm rh} \ge T_{\rm BBN} \simeq 4$ MeV [99–103].

During reheating, the inflaton transfers its energy density to SM radiation, the end of reheating corresponding to the onset of the SM radiation-dominated era. The SM radiation energy density ρ_R is governed by the Boltzmann equation [93]

$$\frac{d\rho_R}{dt} + 4H\,\rho_R = +\frac{2n}{2+n}\,\Gamma_\phi\,\rho_\phi\,. \tag{2.7}$$

Using Eq. (2.5), one can solve Eq. (2.7) and further obtain the solution for the evolution of the radiation energy density during reheating

$$\rho_R(a) \simeq \frac{2\sqrt{3}n}{2+n} \frac{M_P}{a^4} \int_{a_I}^a \Gamma_\phi(a') \sqrt{\rho_\phi(a')} \, a'^3 \, da', \qquad (2.8)$$

where a general scale factor dependence of Γ_{ϕ} has been assumed. Such a dependence can come, for example, from the inflaton mass parameter. The effective mass m_{ϕ} for the inflaton field, understood as the second derivative of its potential, is given by

$$m_{\phi}^{2} \equiv \frac{d^{2}V}{d\phi^{2}} = n\left(n-1\right)\lambda \frac{\phi^{n-2}}{\Lambda^{n-4}} \simeq n\left(n-1\right)\lambda^{\frac{2}{n}} \Lambda^{\frac{2(4-n)}{n}} \rho_{\phi}(a)^{\frac{n-2}{n}}.$$
(2.9)

It is interesting to note that for $n \neq 2$, m_{ϕ} features a field dependence that, in turn, would lead to an inflaton decay rate with a scale factor (or time) dependence.

In the following, two different reheating scenarios will be investigated in which the inflaton either decays to a pair of fermions or bosons via a trilinear coupling.

2.1 Fermionic Reheating

In the case where the inflaton only decays into a pair of fermions Ψ and $\overline{\Psi}$ via a trilinear interaction $y \phi \overline{\Psi} \Psi$, with y being the corresponding Yukawa coupling, the decay rate is given by

$$\Gamma_{\phi} = \frac{y_{\text{eff}}^2}{8\pi} m_{\phi} \,, \tag{2.10}$$

where the effective coupling $y_{\text{eff}} \neq y$ (for $n \neq 2$) is obtained after averaging over oscillations [93, 104, 105]. We will return to our treatment for effective coupling shortly. The evolution of the SM energy density in Eq. (2.8) becomes

$$\rho_R(a) \simeq \frac{\sqrt{3}}{8\pi} \, \frac{n\sqrt{n(n-1)}}{7-n} \, \frac{y_{\text{eff}}^2 \lambda^{\frac{1}{n}} \, M_P \, \rho_\phi(a_{\text{rh}})^{\frac{n-1}{n}}}{\Lambda^{\frac{n-4}{n}}} \left(\frac{a_{\text{rh}}}{a}\right)^{\frac{6(n-1)}{2+n}} \left[1 - \left(\frac{a_I}{a}\right)^{\frac{2(7-n)}{2+n}}\right], \quad (2.11)$$

and, therefore, the temperature of the SM bath evolves as

$$T(a) \simeq T_{\rm rh} \times \begin{cases} \left(\frac{a_{\rm rh}}{a}\right)^{\frac{3}{2}\frac{n-1}{n+2}} & \text{for } n < 7, \\ \frac{a_{\rm rh}}{a} & \text{for } n > 7. \end{cases}$$
(2.12)

The Hubble expansion rate (cf. Eq. (2.6)) can be rewritten as a function of T as

$$H(T) \simeq H_{\rm rh} \begin{cases} \left(\frac{T}{T_{\rm rh}}\right)^{\frac{2n}{n-1}} & \text{for } n < 7, \\ \left(\frac{T}{T_{\rm rh}}\right)^{\frac{3n}{n+2}} & \text{for } n > 7. \end{cases}$$
(2.13)

Before closing this subsection, we note that for the case with n = 2 where the inflaton oscillates in a quadratic potential with an EoS parameter $\omega = 0$, the standard dependences with the scale factor $\rho_R(a) \propto a^{-3/2}$ and $T(a) \propto a^{-3/8}$ are reproduced.

2.2 Bosonic Reheating

Alternatively, if the inflaton only decays into a pair of bosons S through the interaction $\mu \phi S S$, with μ being a coupling with mass dimension, the decay rate is instead

$$\Gamma_{\phi} = \frac{\mu_{\text{eff}}^2}{8\pi \, m_{\phi}} \,, \tag{2.14}$$

where again the effective coupling $\mu_{\text{eff}} \neq \mu$ (if $n \neq 2$) can be obtained after averaging over oscillations. Using a procedure similar to the previous fermionic case, one sees that the SM energy density scales as

$$\rho_R(a) \simeq \frac{\sqrt{3}}{8\pi} \frac{1}{1+2n} \sqrt{\frac{n}{n-1}} \frac{\mu_{\text{eff}}^2}{\lambda^{1/n}} M_P \Lambda^{\frac{n-4}{n}} \rho_\phi(a_{\text{rh}})^{\frac{1}{n}} \left(\frac{a_{\text{rh}}}{a}\right)^{\frac{6}{2+n}} \left[1 - \left(\frac{a_I}{a}\right)^{\frac{2(1+2n)}{2+n}}\right], \quad (2.15)$$

with which the SM temperature evolves as

$$T(a) \simeq T_{\rm rh} \left(\frac{a_{\rm rh}}{a}\right)^{\frac{3}{2}\frac{1}{2+n}},$$
 (2.16)

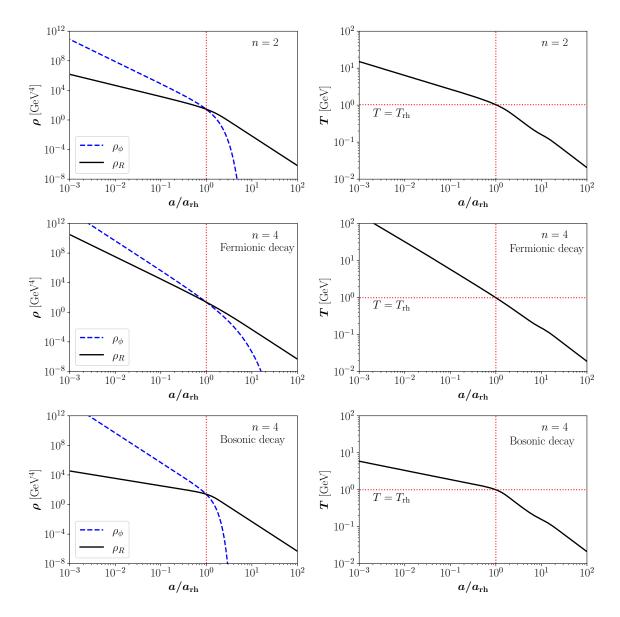


Figure 1. Left: Numerical evolution of the energy densities of the SM (solid black) and ϕ (dashed blue), as a function of the scale factor, for $T_{\rm rh} = 1$ GeV. Right: Corresponding evolution of the SM temperature.

and Hubble as

$$H(T) \simeq H_{\rm rh} \left(\frac{T}{T_{\rm rh}}\right)^{2n},$$
 (2.17)

during reheating.

In Fig. 1, the left panels depict the evolution of both the SM radiation (solid black) and the inflaton (dashed blue) energy densities, as a function of the scale factor, for $T_{\rm rh} = 1$ GeV. The corresponding evolution of the SM bath temperature is shown in the right panels. We note that the bump at $T \sim 0.1$ GeV is due to the QCD phase transition, where the latter leads to a sudden decrease of effective relativistic degrees of freedom in the thermal plasma. The upper panels correspond to the case n = 2, where the effects of bosonic and fermionic decays are equivalent. During reheating, $\rho_{\phi} \propto a^{-3}$, $\rho_R \propto a^{-3/2}$, and $T \propto a^{-3/8}$. The middle and lower panels correspond to n = 4 with fermionic or bosonic decay, respectively. For the first case, $\rho_R \propto a^{-3}$ (and $T \propto a^{-3/4}$) while for the latter case, $\rho_R \propto a^{-1}$ (and $T \propto a^{-1/4}$). Note that to numerically solve the Boltzmann equations for ρ_{ϕ} and ρ_R , the values of the effective couplings appearing in the decay rates are needed, which shall be obtained numerically [93]. Instead, here we fix $T_{\rm rh}$ and then calculate the effective coupling values so that the reheating temperature matches the desired value. Let us remind the reader that the reheating temperature has been defined as $\rho_R(T_{\rm rh}) = \rho_{\phi}(T_{\rm rh}) = 3 M_P^2 H_{\rm rh}^2$.

Before proceeding, a comment on the preheating effect is necessary. With the help of lattice simulations, it has been shown that, during reheating, the EoS parameter $\omega \to 1/3$ for $n \gtrsim 3$ due to inflaton fragmentation and parametric resonance [106–109]. This implies that preheating efficiently drives the Universe to a radiation-dominated phase. However, to fully deplete the inflaton energy, perturbative decay through trilinear couplings between the inflaton and the daughter particles is still required, which is expected to occur in the last stage of the heating process after inflation [107]. Note that for efficient inflaton energy transfer via preheating, sizable couplings are usually required, which could spoil the inflaton potential and inflationary predictions due to loop corrections. Furthermore, in the literature, coherent oscillations of the inflaton condensate have been found to break down [110], which could delay preheating. It has also been shown that depending on the spin of the daughter particles, preheating might be shut down due to the large effective vacuum expectation value of the Higgs field [111]. In this work, we are interested in low-reheating-temperature scenarios, so we assume small inflaton couplings to daughter particles, in which case preheating effects are not expected to be efficient. Thereafter, throughout this work, we will focus on the perturbative reheating scenario.

Having understood possible cosmological histories for the background, in the next section the evolution of the DM number density will be studied, taking particular care to the case where DM freezes out during reheating.

3 Dark Matter Production

The evolution of the DM number density $n_{\rm dm}$ is governed by the Boltzmann equation

$$\frac{dn_{\rm dm}}{dt} + 3H n_{\rm dm} = -\langle \sigma v \rangle \left(n_{\rm dm}^2 - n_{\rm eq}^2 \right), \qquad (3.1)$$

with the Hubble expansion rate

$$H^2 = \frac{\rho_R + \rho_\phi}{3 M_P^2} \,. \tag{3.2}$$

 $n_{\rm eq}(T)$ corresponds to the equilibrium number density given by

$$n_{\rm eq}(T) \simeq g \left(\frac{m\,T}{2\pi}\right)^{3/2} e^{-\frac{m}{T}} \tag{3.3}$$

for non-relativistic particles, where m and g denote the mass and the number of degrees of freedom of the DM field, respectively. In this work, we focus on a temperature-independent cross section, which corresponds to an S-wave dominated interaction. We note that direct production of DM out of the inflaton decay may also be relevant. In that case, Eq. (3.1) would have an additional term $\propto 2 \operatorname{Br} \Gamma_{\phi} n_{\phi}$, where Br denotes the branching ratio of the inflaton decay into a couple of DM particles, and n_{ϕ} is the inflaton number density. Here we focus on the effect of the time dependence of Γ_{ϕ} on the evolution of the WIMP thermal production, and therefore, in the following analysis this contribution will be omitted.³ Finally, for a discussion regarding DM production during the SM thermalization process, we refer the reader to Refs. [112–116].

3.1 Freeze out after Reheating

We first revisit the standard scenario, where DM freeze-out occurs well *after* reheating, in the standard radiation-dominated era. In this case, since the SM entropy is conserved during and after the DM freeze-out, Eq. (3.1) can conveniently be rewritten as

$$\frac{dY}{dx} = -\frac{\langle \sigma v \rangle s}{Hx} \left(Y^2 - Y_{\rm eq}^2 \right), \qquad (3.4)$$

where $x \equiv m/T$, the DM yield $Y \equiv n_{\rm dm}/s$, and the SM entropy density $s(T) = \frac{2\pi^2}{45} g_{\star s} T^3$, with $g_{\star s}(T)$ the number of relativistic degrees of freedom that contribute to the SM entropy. Taking into account that in a radiation-dominated era $H(T) = \sqrt{\frac{\rho_R}{3M_P^2}} = \frac{\pi}{3}\sqrt{\frac{g_{\star}(T)}{10}} \frac{T^2}{M_P}$, with $g_{\star}(T)$ being the number of relativistic degrees of freedom contributing to the SM energy density, Eq. (3.4) admits the standard approximate solution

$$Y_0 \simeq \frac{15}{2\pi g_{\star s}} \sqrt{\frac{g_\star}{10}} \frac{1}{M_P T_{\rm fo} \langle \sigma v \rangle}, \qquad (3.5)$$

where Y_0 corresponds to the DM yield at present, long after freeze-out. To match the whole observed DM relic density, it is required that

$$m Y_0 = \Omega_{\rm dm} h^2 \frac{1}{s_0} \frac{\rho_c}{h^2} \simeq 4.3 \times 10^{-10} \text{ GeV},$$
 (3.6)

with $\rho_c \simeq 1.05 \times 10^{-5} h^2 \text{ GeV/cm}^3$ being the critical energy density, $s_0 \simeq 2.69 \times 10^3 \text{ cm}^{-3}$ the present entropy density [117], and $\Omega_{\rm dm} h^2 \simeq 0.12$ the observed DM relic abundance [2]. The temperature $T_{\rm fo}$ at which the DM freeze-out occurs, or equivalently $x_{\rm fo} \equiv \frac{m}{T_{\rm fo}}$, is

The temperature $T_{\rm fo}$ at which the DM freeze-out occurs, or equivalently $x_{\rm fo} \equiv \frac{m}{T_{\rm fo}}$, is defined by the equality $\frac{n_{\rm eq} \langle \sigma v \rangle}{H}\Big|_{x_{\rm fo}} = 1$, and given by

$$x_{\rm fo} = -\frac{1}{2} \mathcal{W}_{-1} \left[-\frac{8\pi^5}{45} \frac{g_{\star}}{g^2} \frac{1}{\left(M_P \, m \, \langle \sigma v \rangle\right)^2} \right], \tag{3.7}$$

where \mathcal{W}_{-1} is the -1 branch of the Lambert function. The observed DM relic abundance is typically matched for cross-sections at the electroweak scale $\langle \sigma v \rangle = \mathcal{O}(10^{-9}) \text{ GeV}^{-2}$ which corresponds to $\mathcal{O}(10^{-26}) \text{ cm}^3/\text{s}$ and $x_{\text{fo}} \sim 25$, with a small logarithmic dependence on the DM mass [6].

3.2 Freeze out during Reheating

Even if the DM freeze-out is typically assumed to occur well *after* the end of reheating, when the Universe is radiation dominated, this is not guaranteed. Alternatively, DM freeze-out could occur *during* reheating. In this case, the SM entropy is not conserved because of the

³This is typically a good assumption as long as $Br \lesssim 10^{-4} m/(100 \text{ GeV})$ [15, 28].

decay of the inflaton into SM particles. Therefore, instead of Eq. (3.4), it is more convenient to rewrite Eq. (3.1) as

$$\frac{dN}{da} = -\frac{\langle \sigma v \rangle}{H a^4} \left(N^2 - N_{\rm eq}^2 \right), \qquad (3.8)$$

where $N \equiv a^3 n_{\rm dm}$ and $N_{\rm eq} \equiv a^3 n_{\rm eq}$. Using Eq. (2.6) for the Hubble parameter during reheating, Eq. (3.8) can be analytically solved as

$$N(a_{\rm rh}) \simeq \frac{6}{2+n} \frac{H_{\rm rh}}{\langle \sigma v \rangle} a_{\rm rh}^3 \left(\frac{a_{\rm rh}}{a_{\rm fo}}\right)^{\frac{-6}{2+n}}, \qquad (3.9)$$

where $a_{\rm fo}$ and $a_{\rm rh}$ correspond to the scale factor when $T = T_{\rm fo}$ and $T = T_{\rm rh}$, respectively. After reheating, namely, for $a > a_{\rm rh}$, the SM entropy is conserved, and therefore $Y_0 \simeq Y(a_{\rm rh})$, implying

$$Y_{0} = \frac{N(a_{\rm rh})}{s(a_{\rm rh}) a_{\rm rh}^{3}} \simeq \frac{6}{2+n} \frac{45}{2\pi^{2} g_{\star s}} \frac{H_{\rm rh}}{\langle \sigma v \rangle T_{\rm rh}^{3}} \times \begin{cases} \left(\frac{T_{\rm rh}}{T_{\rm fo}}\right)^{\frac{4}{n-1}} & \text{for fermionic decay,} \\ \left(\frac{T_{\rm rh}}{T_{\rm fo}}\right)^{4} & \text{for bosonic decay,} \end{cases}$$
(3.10)

for n < 7 in the case of fermionic decay. This assumption will be followed from now on. We note that, in the era after the freeze-out and before the end of reheating, the DM yield evolves as $Y(T) \propto (aT)^{-3}$, which corresponds to

$$Y(T) \propto \begin{cases} T^{\frac{7-n}{n-1}} & \text{for fermionic decay,} \\ T^{1+2n} & \text{for bosonic decay.} \end{cases}$$
(3.11)

Examples of the evolution of DM abundance are shown in Fig. 2, for reference points that fit the total observed abundance: the upper panel corresponds to n = 2, $T_{\rm rh} = 1$ GeV, m = 100 GeV and $\langle \sigma v \rangle = 10^{-11}$ GeV⁻², the lower left panel corresponds to a fermionic decay with n = 4, $T_{\rm rh} = 0.1$ GeV, m = 100 GeV and $\langle \sigma v \rangle = 6 \times 10^{-10}$ GeV⁻², and the lower right panel corresponds to a bosonic decay with n = 4, $T_{\rm rh} = 1$ GeV, m = 40 GeV and $\langle \sigma v \rangle = 5 \times 10^{-11}$ GeV⁻². The solid black lines depict the evolution of the DM yield Y, while the dashed black lines show the equilibrium yield $Y_{\rm eq}$. Furthermore, the red dotted vertical lines show $T = T_{\rm fo}$ and $T = T_{\rm rh}$, and the gray horizontal bands correspond to the observed DM abundance $\Omega_{\rm dm}h^2 \simeq 0.12$ [2].

The evolution of the DM yield depicted in Fig. 2 is obtained by numerically solving the Boltzmann equations (2.4), (2.7), (3.1) and the Friedmann equation (3.2). We have checked that the analytical estimate of the asymptotic value of Y_0 shown in Eq. (3.10) agrees well with the numerical result. Furthermore, the evolution of the DM yield for $T_{\rm fo} \gg T \gg T_{\rm rh}$ (that is, $Y \propto x^{-5}$ for n = 2, and $Y \propto x^{-1}$ or $Y \propto x^{-9}$ for fermionic or bosonic decays, respectively, and n = 4) also matches the analytical expressions of Eq. (3.11). It is worth noticing that in the lower right panel, the DM only reaches chemical equilibrium with the SM bath at $x \simeq 3$.

The contours for the reheating temperature required to match the whole measured DM abundance are shown in Fig. 3, in the parameter space $[m, \langle \sigma v \rangle]$, for bosonic and fermionic decays of the inflaton, and different values of n. The thick black lines correspond to the standard case, where DM freezes out in a radiation-dominated era, whereas the thin black lines to $T_{\rm rh} = T_{\rm BBN}$, 10^{-1} GeV, 10^{0} GeV and 10^{1} GeV. Some constraints apply and are

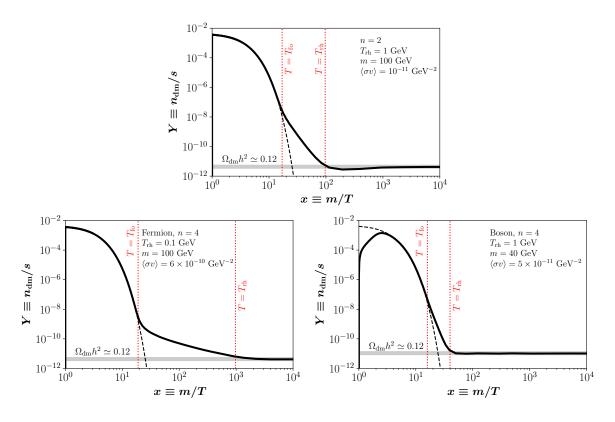


Figure 2. Examples of the evolution of the DM abundance, for different benchmark points described in the text that fit the total observed abundance. The solid black lines show the DM yield Y, whereas the dashed black lines show the equilibrium density. Furthermore, the dotted red vertical lines correspond to $T = T_{\rm fo}$ and $T = T_{\rm rh}$, and the gray horizontal bands depict the observed DM abundance $\Omega_{\rm dm}h^2 \simeq 0.12$.

shown as red regions: *i*) Above the thick black line, corresponding to higher cross sections, DM decouples very late and is therefore underproduced ($\Omega_{\rm dm}h^2 < 0.12$). *ii*) Reheating temperatures below $T_{\rm BBN}$ are in conflict with cosmological observations, which corresponds to $m \leq \mathcal{O}(10^{-1})$ GeV. And *iii*), very small thermally averaged cross sections $\langle \sigma v \rangle$ could not be enough to guarantee chemical equilibrium between the dark and visible sectors. It is worth mentioning that in this case, labeled 'No freeze-out', the whole DM relic abundance could be matched, however, the production would not correspond to the WIMP mechanism, but rather to a FIMP scenario, with the usual strong dependence on initial conditions. This case will not be considered in the present analysis.

In Fig. 3, it can be seen that a DM freeze-out during reheating allows exploring smaller cross-sections $\langle \sigma v \rangle$ compared to the usual case where the freeze-out occurs in the standard radiation-dominated case. For example, with n = 2 and $T_{\rm rh} = 10$ GeV, $\langle \sigma v \rangle$ can be as small as $\mathcal{O}(10^{-15})$ GeV⁻² for a DM mass of ~ 7 TeV. In a scenario with larger $T_{\rm rh}$, even smaller $\langle \sigma v \rangle$ are allowed for heavier WIMPs. We also note that for a fixed $T_{\rm rh}$, $\langle \sigma v \rangle$ decreases with increasing m. This is because for heavier WIMPs, freeze-out occurs earlier (i.e., at a larger $T_{\rm fo}$), which implies a smaller $\langle \sigma v \rangle$. Finally, it should be noted that with increasing n, the parameter space shrinks, as depicted in Fig. 3. The main reason is that the dilution effect becomes less prominent with larger n (that is, the factor $(a_{\rm rh}/a_{\rm fo})^{\frac{-6}{2+n}}$ in Eq. (3.9)), and hence a larger $\langle \sigma v \rangle$ is needed to compensate for this. A larger $\langle \sigma v \rangle$ corresponds to a later freeze-out,

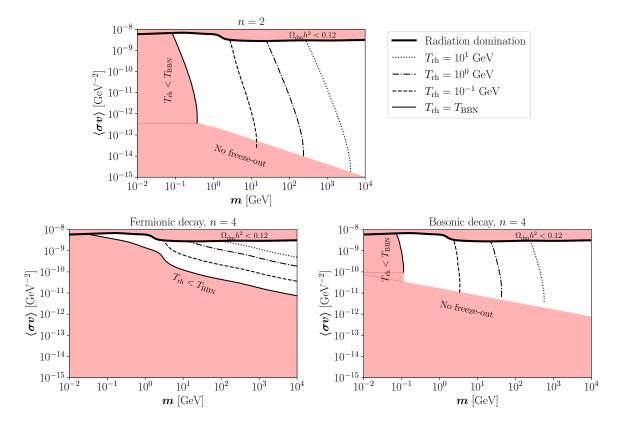


Figure 3. Parameter space that matches the whole observed DM abundance, for $T_{\rm rh} = T_{\rm BBN}$ (solid lines), 10^{-1} GeV (dashed), 10^{0} GeV (dash-dotted), and 10^{1} GeV (dotted). The upper panel corresponds to n = 2, while the lower to n = 4 and fermionic (left) or bosonic (right) decays of the inflaton. The thick black lines correspond to the standard scenario, where DM freeze-out in the radiation-dominated era. The red areas are disregarded as they generate a DM underabundance ($\Omega_{\rm dm} < 0.12$), the reheating occurs after BBN ($T_{\rm rh} < T_{\rm BBN}$), or DM does not reach chemical equilibrium with the SM ('No freeze-out').

where the WIMP densities are smaller, and therefore a smaller dilution is required. In the present case, DM freezes out during reheating, at a temperature given by

$$x_{\rm fo} \simeq \begin{cases} \frac{3+n}{2(1-n)} \mathcal{W}_{-1} \left[\frac{2(1-n)}{3+n} \left(\frac{g^2}{(2\pi)^3} \frac{m^2 T_{\rm rh}^4 \langle \sigma v \rangle^2}{H_{\rm rh}^2} \right)^{\frac{1-n}{3+n}} \left(\frac{m}{T_{\rm rh}} \right)^{\frac{4}{3+n}} \right] & \text{for fermions,} \\ \frac{3-4n}{2} \mathcal{W}_{-1} \left[\frac{2}{3-4n} \left(\frac{g^2}{(2\pi)^3} \frac{m^2 T_{\rm rh}^4 \langle \sigma v \rangle^2}{H_{\rm rh}^2} \right)^{\frac{1}{3-4n}} \left(\frac{m}{T_{\rm rh}} \right)^{\frac{4(1-n)}{3-4n}} \right] & \text{for bosons.} \end{cases}$$
(3.12)

Figure 4 shows the freeze-out temperature required to match the entire observed DM abundance, as a function of the DM mass, for $T_{\rm rh} = T_{\rm BBN}$ (solid lines), 10^{-1} GeV (dashed lines), 10^0 GeV (dash-dotted lines) and 10^1 GeV (dotted lines). The upper panel corresponds to n = 2, the lower left to a fermionic decay with n = 4, and the lower right to a bosonic decay with n = 4. The thick black lines represent the standard freeze-out in radiation domination; cf. Eq. (3.7). Additionally, the lower red bands depict a relativistic freeze-out, i.e. $x_{\rm fo} \leq 3$. Note that for n = 2, the freeze-out temperature can be as low as $x_{\rm fo} \sim 6$ due to the entropy

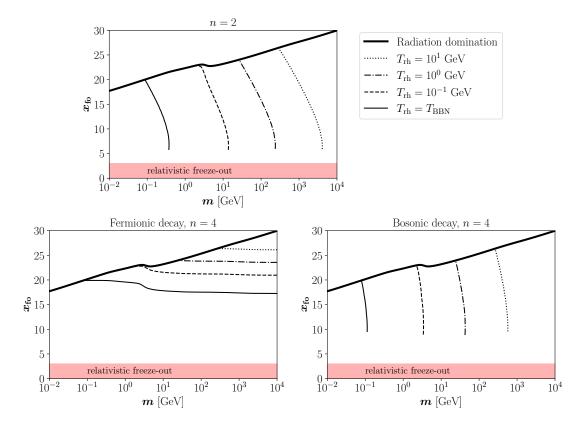


Figure 4. Freeze-out temperature required to match the whole observed DM abundance, as a function of the DM mass, for $T_{\rm rh} = T_{\rm BBN}$ (solid lines), 10^{-1} GeV (dashed lines), 10^{0} GeV (dotted lines in dashes) and 10^{1} GeV (dotted lines). The upper panel corresponds to n = 2, while the lower left to a fermionic decay with n = 4, and the lower right to a bosonic decay with n = 4. The thick black lines represent the standard freeze-out in radiation domination.

dilution effect. For n = 2 and n = 4 (bosonic), the lines are cut to ensure that WIMPs thermalize; lower values of x_{fo} would correspond to the 'No freeze-out' regime depicted in Fig. 3. For a fermionic decay with n = 4, x_{fo} tends to be independent of m, with only a small dependence on the relativistic degrees of freedom. Note that the temperature during reheating (for the fermionic case) tends to feature the same scaling as the free radiation for large values of n, cf. Eq. (2.12); this implies that the corresponding freeze-out behavior becomes similar to the standard case. Actually, such a tendency can also be seen in Fig. 4. However, for the bosonic case, due to a different temperature scaling (cf. Eq. (2.16)), the allowed parameter space becomes distinct from the fermionic case with an increase of n. Note that, as argued earlier, in both fermonic and bosonic cases, with increasing n, a larger $\langle \sigma v \rangle$ (corresponding to a later freeze-out with lower T_{fo}) is needed, implying that x_{fo} increases with n, as depicted in Fig. 4.

Before closing this section, we comment on some present and future possibilities of testing the current scenario using the different channels offered by DM indirect detection experiments. In addition to the information given in Fig. 3, Fig. 5 also shows in blue present constraints coming from: *i*) CMB spectral distortions from WIMP annihilation into charged particles [118] (labeled 'CMB'), *ii*) the combined analysis (labeled 'combined') of γ -ray data using Fermi-LAT, HAWC, H.E.S.S., MAGIC, and VERITAS by considering WIMP annihilation to $\tau^+\tau^-$

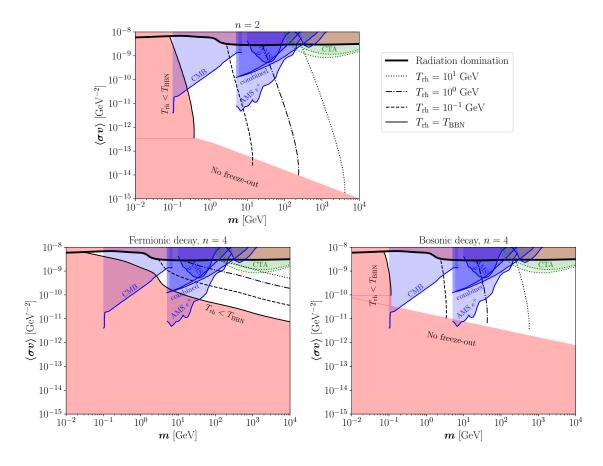


Figure 5. Same as Fig. 3 with current experimental bounds (blue) and future sensitivities (green).

(upper line) and $b\bar{b}$ (lower line) [119], and *iii*) AMS-02 measurements of antimatter in cosmic rays, in the positron [120] and the antiproton [121] channels.⁴ Additionally, Fig. 5 shows in green the projected sensitivity of the ground-based CTA experiment, which could be capable of testing WIMP DM at the TeV scale and above. The upper (lower) green dotted line shows its expected sensitivity for WIMP annihilation to $b\bar{b}$ (W^+W^-) [129]. Note that CTA might be able to probe a large fraction of the TeV-scale WIMPs in the standard scenario, and even a chunk of the parameter space where the freeze-out occurs during reheating, for both fermionic and bosonic decays.

4 Conclusions

Despite the huge experimental efforts over the past decades, the nature of dark matter (DM) remains one of the most challenging and fundamental questions in particle physics. In particular, scenarios where DM is a weakly interacting massive particle (WIMP) reaching thermal equilibrium with the standard model (SM) bath have received immense attention both theoretically and experimentally, but unfortunately no overwhelming evidence for WIMP DM has been found. This motivates a quest beyond the standard WIMP paradigm. An alternative

⁴It is interesting to recall that the bounds presented depends on the DM annihilation channels, and suffer from large uncertainties arising from e.g. assumptions on the DM density profile, the local DM density, and the propagation of charged particles in the interstellar medium [122–128].

is to renounce the thermal nature of DM, as in the FIMP paradigm. However, these kinds of scenario suffer from a strong dependence on the largely unknown initial conditions after cosmological inflation.

In this work, we focus on another avenue where DM remains a thermal relic, but its freeze-out (that is, its departure from chemical equilibrium) occurs during (and not after) inflationary reheating. In particular, we investigated a scenario where the inflaton field ϕ coherently oscillates at the bottom of a generic potential $V(\phi) \propto \phi^n$, while decaying into SM particles to reheat the Universe. Depending on the details of the reheating (i.e. the shape of the potential and the coupling between the inflaton and the SM particles), the inflaton and the SM energy density could feature a distinct evolution. Consequently, in this case the behavior of WIMP freeze-out differs from the case where it occurs well after reheating. Due to the dilution effect during reheating, smaller thermally averaged cross-sections $\langle \sigma v \rangle$ becomes compatible with the measured DM relic abundance.

In particular, if DM freezes-out after reheating, one requires $\langle \sigma v \rangle \sim \mathcal{O}(10^{-9}) \text{ GeV}^{-2} \sim \mathcal{O}(10^{-26}) \text{ cm}^3/\text{s}$ to match the observed DM abundance, and a DM freeze-out temperature $x_{\text{fo}} \equiv m/T_{\text{fo}} \sim 25$. Instead, if the freeze-out occurs during reheating while the inflaton oscillates in a quadratic potential (n = 2), if $T_{\text{rh}} \sim 1$ GeV, $\langle \sigma v \rangle$ could be of the order $\mathcal{O}(10^{-15}) \text{ GeV}^{-2}$ for a WIMP at the TeV scale. Interestingly, a sizable part of the parameter space is in the range of current experiments and future proposals. For instance, for TeV-scale WIMPs, the upcoming CTA could probe the expected cross-sections for DM decoupling after reheating, and a portion of the plane $[m, \langle \sigma v \rangle]$ relevant for the case where DM decouples during reheating.

Acknowledgments

We would like to thank Manuel Drees once again for enlightening discussions. NB received funding from the Patrimonio Autónomo - Fondo Nacional de Financiamiento para la Ciencia, la Tecnología y la Innovación Francisco José de Caldas (MinCiencias - Colombia) grants 80740-465-2020 and 80740-492-2021. NB is also funded by the Spanish FEDER/MCIU-AEI under grant FPA2017-84543-P. This project has received funding/support from the European Union's Horizon 2020 research and innovation program under the Marie Skłodowska-Curie grant agreement No 860881-HIDDeN.

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