Self-interacting Dark Matter from Primordial Black Holes

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Abstract. The evaporation of primordial black holes (PBH) with masses ranging from \(\sim 10^{-1}\) to \(\sim 10^9\) g could have generated the whole observed dark matter (DM) relic density. It is typically assumed that after being produced, its abundance freezes and remains constant. However, thermalization and number-changing processes in the dark sector can have a strong impact, in particular enhancing the DM population by several orders of magnitude. Here we estimate the boost from general arguments such as the conservation of energy and entropy, independently from the underlying particle physics details of the dark sector. Two main consequences can be highlighted: \textit{i}) As the DM abundance is increased, a smaller initial energy density of PBHs is required. \textit{ii}) Thermalization in the dark sector decreases the mean DM kinetic energy, relaxing the bound from structure formation and hence, allowing light DM with mass in the keV ballpark.
1 Introduction

The existence of dark matter (DM) has been firmly established by astrophysical and cosmological observations, although its fundamental nature remains elusive [1]. Up to now, the only evidence about the existence of such a dark component is via its gravitational interactions. In the last decades, weakly interacting massive particles (WIMPs), with masses and couplings at the electroweak scale, have been the leading DM production paradigm [2]. However, the increasingly strong observational constraints on DM are urging the quest for alternative scenarios.

Several alternatives to the classical WIMP mechanism exist. For instance, one can have deviations from the standard expansion history of the early universe [3]. Another possibility occurs if the couplings between the dark and visible sectors are very suppressed, so that DM never reaches chemical equilibrium with the standard model (SM), as in the case of the so-called freeze-in mechanism (FIMP) [4–8]. An extreme case occurs if DM is only coupled to the SM through Planck suppressed higher dimensional operators, and is produced via purely gravitational interactions [9–13].

However, the DM genesis could also be intimately related to the Hawking evaporation of primordial black holes (PBH). In fact, PBHs could have been formed from inhomogeneities in the early universe [14]. If their initial mass was below ~ 10^9 g, they disappear through Hawking evaporation [15] before Big Bang nucleosynthesis (BBN), and are poorly constrained [16, 17]. During the evaporation, PBHs radiate not only SM particles but also hidden sector states, and in particular DM. In this regard, PBH evaporation may have played a central role in the DM production [18–28].
Previous studies typically assumed that after production via PBH evaporation, the DM abundance remains constant. Nevertheless, even if there is no DM production out of the visible sector, the dynamics in the dark sector could be not trivial, featuring for example $N$-to-$N'$ number-changing interactions, where $N$ DM particles annihilate into $N'$ of them (with $N > N' \geq 2$). The dominant $N$-to-$N'$ processes are naturally 3-to-2 (see e.g. Refs. [29–37]), but are forbidden in models where DM is protected by a $Z_2$ symmetry. In that case, unavoidable 4-to-2 annihilations [38–43] could dominate. If these processes reach equilibrium, DM forms a thermal bath with a temperature in general different from the one of the SM. More importantly, number-changing processes have a strong impact on DM, increasing by several orders of magnitude its relic abundance [33, 38, 40, 44–50].

In this work, we investigate the impact of thermalization of the dark sector on the DM abundance produced by evaporation of PBHs. For that purpose, in section 2 we briefly revisit the formation and evaporation of PBH, whereas in section 3 the standard DM production via Hawking radiation is presented. Section 4 is devoted to quantify the DM self-interaction effects on its abundance, and contains our main results. Finally, in section 5 our conclusions are presented.

2 The Rise and Fall of PBHs

Formation and evaporation of PBHs has been vastly discussed in the literature, see, for instance, Refs. [16, 17, 26, 28]. Here we briefly review the main aspects.

2.1 Formation

PBH formed in a radiation dominated epoch, when the SM plasma has a temperature $T = T_{in}$, have an initial mass $M_{in}$ similar to the enclosed mass in the particle horizon, given by

$$M_{in} \equiv M_{BH}(T_{in}) = \frac{4\pi}{3} \, \gamma \, \rho_R(T_{in}) \, H^3(T_{in}).$$

In this expression, \( \gamma \approx w^{3/2} \approx \frac{1}{3} \) (the equation of state parameter \( w = \frac{1}{3} \) in a radiation-dominated epoch), $\rho_R(T) = \frac{\pi^2}{30} \, g_*(T) \, T^4$ is the SM radiation energy density with $g_*(T)$ the number of relativistic degrees of freedom contributing to $\rho_R$ [51], and $H^2(T) = \frac{\rho(T)}{3M^2_P}$ is the squared Hubble expansion rate in terms of the total energy density $\rho(T)$, with $M_P$ the reduced Planck mass.

Extended PBH mass functions arise naturally if the PBHs are created from inflationary fluctuations or cosmological phase transitions, see e.g. Refs. [16, 17, 52–58]. However, for the sake of simplicity, in the present analysis we assume that all the PBHs have the same mass (i.e. they are produced at the same temperature), which is a usual assumption in the literature. Finally, PBHs can gain mass via mergers [25, 59, 60] and accretion [26, 61, 62]. These processes are typically not very efficient, inducing a mass gain of order $O(1)$, and will be hereafter ignored.

2.2 Evaporation

PBHs evaporate by emitting particles lighter than its temperature $T_{BH}$ via Hawking radiation [15]. Given the fact that Hawking radiation can be described as blackbody radiation (up

\[\text{See however Ref. [28], were scenarios with a second DM production mechanism have been considered.}\]
to greybody factors), the energy spectrum of a species \( j \) with \( g_j \) internal degrees of freedom radiated by a nonrotating BH with zero charge is therefore [28, 63]

\[
\frac{d^2 u_j(E, t)}{dt \, dE} = \frac{g_j}{8\pi^2} \frac{E^3}{e^{E/T_{BH}} + 1}, \quad (+ \text{ for fermions, } - \text{ for bosons})
\]

(2.2)

where \( u_j \) is the total radiated energy per unit area, \( t \) the time, \( E \) the energy of the emitted particle \( j \), and

\[
T_{BH} = \frac{M_B^2}{M_{BH}} \simeq 10^{13} \text{GeV} \left( \frac{1}{M_{BH}} \right)
\]

(2.3)

the BH horizon temperature.

The evolution of the BH mass due to Hawking evaporation is given by

\[
\frac{dM_{BH}}{dt} = -4\pi r_S^2 \sum_j \int_0^\infty \frac{d^2 u_j(E, t)}{dt \, dE} dE = -\frac{\pi g_\star(T_{BH})}{480} \frac{M_B^2}{M_{BH}^4},
\]

(2.4)

where \( r_S \equiv \frac{M_{BH}}{4\pi M_P^2} \) is the Schwarzschild radius of the BH. Assuming that \( g_\star \) has no temperature dependence during the whole lifetime of the BH, Eq. (2.4) admits the analytical solution

\[
M_{BH}(t) = M_{in} \left( 1 - \frac{t - t_{in}}{\tau} \right)^{1/3},
\]

(2.5)

where \( t_{in} \) corresponds to the time at formation of the PBH and

\[
\tau \equiv \frac{160}{\pi g_\star(T_{BH})} \frac{M_{in}^3}{M_P^4}
\]

(2.6)

is the PBH lifetime.

Complete BH evaporation happens at \( t = t_{ev} \equiv t_{in} + \tau \simeq \tau \) at which \( M_{BH}(t_{ev}) = 0 \).

In a universe dominated by radiation during the whole BH lifetime, it corresponds to a temperature

\[
T_{ev} \equiv T(t_{ev}) \simeq \left( \frac{9 g_\star(T_{BH})}{10240} \right)^{1/3} \left( \frac{M_P^5}{M_{in}^3} \right)^{1/8} \simeq 1.2 \times 10^{10} \text{ GeV} \left( \frac{g_\star(T_{BH})}{106.75} \right)^{1/3} \left( \frac{1}{M_{in}} \right)^{3/2},
\]

(2.7)

using the fact that \( H(t) = 1/(2t) \) in a radiation-dominated era.

Additionally, the total number \( N_j \) of the species \( j \) of mass \( m_j \) emitted during the PBH evaporation is

\[
N_j = \int_{t(m_j)}^{t_{ev}} \int_0^\infty dE \frac{d^2 N_j}{dt \, dE} = \int_{t(m_j)}^{t_{ev}} dt \int_0^\infty dE \frac{4\pi r_S^2}{E^2} \frac{d^2 u_j}{dt \, dE},
\]

(2.8)

where \( t(m_j) \) corresponds to the time at which BH start emitting \( j \) particles, i.e. when \( T_{BH} \geq m_j \), and is given by

\[
t(m_j) = \max \left[ t_{in}, \ t_{in} + \tau \left( 1 - \left[ \frac{M_B^2}{m_j M_{BH}} \right]^{3/2} \right) \right].
\]

(2.9)

\footnote{Where it has been taken into account that \( \frac{t_{in}}{\tau} \approx 7.86 \times 10^{-12} \left( \frac{g_\star(T_{BH})}{106.75} \right) \left( \frac{0.2}{\gamma} \right) \left( \frac{1}{M_{in}} \right)^{3/2} \ll 1. \)
Therefore, Eq. (2.8) reduces to

\[
N_j = \frac{15 \zeta(3)}{\pi^4} \frac{g_j C_n}{g_*(T_{BH})} \left\{ \begin{array}{ll}
\left( \frac{M_p}{M_j} \right)^2 & \text{for } m_j \leq T_{BH}^n, \\
\left( \frac{M_p}{m_j} \right)^2 & \text{for } m_j > T_{BH}^n,
\end{array} \right.
\]

(2.10)

where \( T_{BH}^n \equiv T_{BH}(t = t_{in}) \) is the initial PBH temperature, and \( C_n = 1 \) or \( 3/4 \) for bosonic or fermionic species, respectively. Additionally, their mean energy is \([21, 24, 64]\)

\[
\langle E_j \rangle = \frac{1}{N_j} \int_{t(m_j)}^{t_{ev}} dt \int_0^\infty dE 3T_{BH} \frac{d^2N_j}{dt dE} = \left\{ \begin{array}{ll}
6T_{BH}^n & \text{for } m_j \leq T_{BH}^n, \\
6m_{DM} & \text{for } m_j > T_{BH}^n,
\end{array} \right.
\]

(2.11)

where \( 3T_{BH} \) is the average energy of particles radiated by a PBH with temperature \( T_{BH} \).\(^3\)

PBH evaporation produces all particles, and in particular extra radiation that can modify successful BBN predictions. To avoid it, we require PBHs to fully evaporate before BBN time, i.e. \( T_{ev} > T_{BBN} \simeq 4 \text{ MeV} \ [65–69] \), which translates into an upper bound on the initial PBH mass\(^4\)

\[
M_{in} \lesssim 2 \times 10^8 \text{ g}.
\]

(2.12)

On the opposite side, a lower bound on \( M_{in} \) can be set once the upper bound on the inflationary scale is taken into account. The limit reported by the Planck collaboration \( H_I \lesssim 2.5 \times 10^{-5}\,M_P \ [70] \) implies that

\[
M_{in} \gtrsim 4\pi \gamma \frac{M_P^2}{H_I} \simeq 0.1 \text{ g}.
\]

(2.13)

Before concluding this section, let us note that as BHs scale like non-relativistic matter \( (\rho_{BH} \propto a^{-3}) \), its energy density \( \rho_{BH} \) naturally tends to dominate over the SM energy density that scales like \( \rho_R \propto a^{-4} \). The initial PBH energy density is usually normalized to the SM energy density at the time of formation \( T = T_{in} \) via the dimensionless parameter

\[
\beta \equiv \frac{\rho_{BH}(T_{in})}{\rho_R(T_{in})} = \frac{M_{in}n_{in}}{\rho_R(T_{in})},
\]

(2.14)

where \( n_{in} \) is the initial BH number density. A matter-dominated era (i.e. a PBH domination) can be avoided if \( \rho_{BH} \ll \rho_R \) at all times, or equivalently if

\[
\beta \ll \frac{T_{ev}}{T_{in}}.
\]

(2.15)

### 3 Dark Matter Production

The whole observed DM relic abundance could have been Hawking radiated by PBHs.\(^5\) The DM production can be analytically computed in two limiting regimes where PBHs dominated or not the energy density of the universe, and will be presented in the following.

\(^3\)The emission is not exactly blackbody but depends upon the spin and charge of the emitted particle [63].

\(^4\)The corresponding bound for BH domination decreases by a factor \((3/4)^{1/3}\).

\(^5\)For the sake of completeness, notice that even if the \( s \)-channel exchange of a graviton gives an irreducible contribution to the total DM relic abundance [9–12], it will be hereafter disregarded.
3.1 Radiation Dominated Universe

The DM yield $Y_{\text{DM}}$ is defined as the ratio of the DM number density $n_{\text{DM}}$ and the SM entropy density $s(T) = \frac{2\pi^2}{45}g_{\ast\ast}(T)T^3$, where $g_{\ast\ast}(T)$ is the number of relativistic degrees of freedom contributing to the SM entropy [51].

In a radiation dominated universe, the DM yield produced by Hawking evaporation of PBHs can be estimated by

$$Y_{\text{DM}} \equiv \frac{n_{\text{DM}}(T_0)}{s(T_0)} = N_{\text{DM}} \frac{n_{\text{in}}}{s(T_{\text{in}})} = \frac{3}{4} N_{\text{DM}} \frac{g_{\ast}(T_{\text{BH}})}{g_{\ast\ast}(T_{\text{BH}})} \frac{T_{\text{in}}}{M_{\text{in}}},$$

with $T_0$ the SM temperature at present, and where the conservation of SM entropy was used. Additionally, $N_{\text{DM}}$ is total number of DM particles emitted by a PBH, and is given by Eq. (2.10).

3.2 Matter Dominated Universe

Alternatively, PBHs can dominate the universe energy density before their decay. In that case, the DM yield is instead

$$Y_{\text{DM}} \equiv \frac{n_{\text{DM}}(T_0)}{s(T_0)} = \frac{n_{\text{DM}}(T_{\text{ev}})}{s(T_{\text{ev}})} = N_{\text{DM}} \frac{n_{\text{BH}}(t_{\text{ev}})}{s(T_{\text{ev}})},$$

using again the conservation of the SM entropy after the PBHs have completely evaporated, and where $T_{\text{ev}}$ is the SM temperature just after the complete BH evaporation. Additionally, assuming an instantaneous evaporation of the BHs at $t = t_{\text{ev}} \approx \tau$, one has that

$$n_{\text{BH}}(t_{\text{ev}}) = \frac{\rho_{\text{BH}}(t_{\text{ev}})}{M_{\text{in}}} = \frac{3M_P^2H^2(t_{\text{ev}})}{M_{\text{in}}} \approx \frac{4M_P^2}{3M_{\text{in}}t_{\text{ev}}^2} \simeq \frac{\pi^2g_{\ast}^4(T_{\text{BH}})}{19200} \frac{M_{P}^{10}}{M_{\text{in}}^2},$$

and that

$$T_{\text{ev}}^4 \simeq \frac{g_{\ast}(T_{\text{BH}})}{640} \frac{M_{P}^{10}}{M_{\text{in}}},$$

where the fact that in a matter-dominated universe $H(t) = 2/(3t)$ was used. Therefore, the DM yield in Eq. (3.2) can be expressed as

$$Y_{\text{DM}} \simeq \frac{3}{4} N_{\text{DM}} \frac{g_{\ast}(T_{\text{BH}})}{g_{\ast\ast}(T_{\text{BH}})} \frac{T_{\text{ev}}}{M_{\text{in}}}. \quad (3.5)$$

We notice that, as expected, the DM yields in the radiation dominated (Eq. (3.1)) and matter dominated (Eq. (3.5)) eras become identical in the limit $\beta \to T_{\text{ev}}/T_{\text{in}}$, cf. Eq. (2.15), with $T_{\text{ev}} = T_{\text{ev}}$.

To reproduce the observed DM relic abundance $\Omega_{\text{DM}}h^2 \simeq 0.12$ [1], the DM yield has to be fixed so that $m_{\text{DM}}Y_{\text{DM}} = \Omega_{\text{DM}}h^2 \frac{1}{s_0} \frac{\rho_c}{h^2} \simeq 4.3 \times 10^{-10}$ GeV, where $\rho_c \simeq 1.1 \times 10^{-5} h^2$ GeV/cm$^3$ is the critical energy density, and $s_0 \simeq 2.9 \times 10^3$ cm$^{-3}$ is the entropy density at present [1]. Figure 1 shows with thick black lines the parameter space reproducing the observed DM density for different DM masses. The shaded regions represent areas constrained by different observables: $M_{\text{in}} \lesssim 10^{-1}$ g and $M_{\text{in}} \gtrsim 2 \times 10^8$ g are disfavored by

\footnote{In the approximation of an instantaneous evaporation of the PBHs, the SM entropy density is violated at $t = t_{\text{ev}}$, and therefore there is a sudden increase of the SM temperature, from $T_{\text{ev}}$ to $T_{\text{ev}}$.}
Figure 1. Parameter space reproducing the observed DM abundance (thick black lines) from PBH evaporation, \textit{without} DM self-interactions. The shaded areas are excluded by different observables described in the text. The hot DM bound \textit{only} applies to $m_{\text{DM}} \lesssim 30$ GeV.

CMB and BBN (both in red), $\beta$ values smaller than $\sim 10^{-24}$ can not accommodate the total observed DM abundance (green) (Eq. (3.6)), whereas large values for $\beta \gtrsim 10^{-5}$ produce hot DM (blue) (Eq. (3.9)). It is important to note that the latter constraint \textit{only} applies to the case of light DM ($m_{\text{DM}} \ll T_{\text{BH}}$), with mass $m_{\text{DM}} \lesssim 30$ GeV. Finally, the dotted red line shows the transition between radiation (lower part) and matter-dominated eras (upper part).

In this figure, the effects of both the SM-DM interactions and the DM self-interactions have been neglected, and therefore it corresponds to DM produced solely via Hawking evaporation of PBHs. The thick black lines show three different slopes, corresponding to three different regimes. If PBHs dominated the universe energy density (above the red dotted line), the DM yield is independent of $\beta$, cf. Eq. (3.5), and therefore the lines are vertical. In this regime, $m_{\text{DM}} \simeq 10^9$ GeV is the lowest viable DM mass. However, a $\beta$ dependence shows up if, during the whole BH lifetime, the universe was radiation dominated (below the dotted line), Eq. (3.1). In this case, two regimes arise, depending on whether DM is lighter or heavier than the initial BH temperature, Eq. (2.10). In the former case $\beta \propto T_{\text{in}}$, whereas in the latter $\beta \propto T_{\text{in}}^{-3}$.

In the present case where DM self-interactions are not efficient, DM has to be heavier than $\mathcal{O}(1)$ MeV in order not to be hot, and can be as heavy as $M_P$ [71, 72] (notice that we are not considering a BH evaporation process stopping at $T_{\text{BH}} \sim M_P$, with the associated production of Planck mass relics [25, 64, 73–76]).

Before closing the section, two comments are in order. On the one hand, as mentioned previously, if one requires the PBHs to radiate the whole observed DM abundance, a lower
limit on $\beta$ appears in the radiation-dominated era when $m_{DM} = T_{in}^{BH}$, and corresponds to
\[ \beta \geq \frac{4\pi^4}{45\zeta(3)} C_n \frac{g_{s}(T_{in})}{g_{DM}} \frac{m_{DM} Y_{DM}}{T_{in}} \simeq 7.6 \times 10^{-23} \left( \frac{g_{s}(T_{BH})}{106.75} \right)^{5/4} \left( \frac{0.2}{\gamma} \right)^{1/2} \left( \frac{M_{in}}{1 \text{ g}} \right)^{1/2}. \] (3.6)

This bound is shown in green in Fig. 1. On the other hand, we note that due to their large initial momentum, DM particles could have a large free-streaming length leading to a suppression on the structure formation at small scales. In the present scenario where DM has no interactions with the SM or with itself, the DM momentum simply redshifts, and its value $p_0$ at the present time is [21]
\[ p_0 = \frac{a_{ev}}{a_0} p_{ev} = \frac{a_{ev}}{a_{eq}} \frac{m_{DM}}{a_0} p_{ev} = \frac{a_{ev}}{a_{eq}} \frac{\Omega_R}{\Omega_m} p_{ev} = \left[ \frac{g_{s}(T_{eq})}{g_{ss}(T_{ev})} \right]^{1/3} \frac{T_{eq}}{T_{ev}} \frac{\Omega_R}{\Omega_m} p_{ev}, \] (3.7)

where $T_{eq}$ and $a_{eq}$ correspond to the temperature and the scale factor at the matter-radiation equality, respectively. For light DM ($m_{DM} \ll T_{BH}^{in}$), $p_{ev} \simeq T_{BH}^{in}$ and by using Eq. (2.7), the DM typical momentum at the present time can be estimated as
\[ p_0 \simeq \left[ \frac{g_{s}(T_{eq})}{g_{ss}(T_{ev})} \right]^{1/3} \left[ \frac{10240}{9 g_{s}(T_{BH})} \right]^{1/4} \frac{T_{ev}}{T_{eq}} \frac{\Omega_R}{\Omega_m} \left[ \frac{M_{in}}{M_{P}} \right]^{1/2} \simeq 7.8 \times 10^{-14} \left[ \frac{M_{in}}{M_{P}} \right]^{1/2} \text{GeV}, \] (3.8)

where $T_{eq} \simeq 0.8 \text{ eV}$, $\Omega_R \simeq 5.4 \times 10^{-5}$ and $\Omega_m \simeq 0.315$ [1, 77] were used. A lower bound on the DM mass can be obtained from the upper bound on a typical velocity of warm DM at present time. Taking $v_{DM} \lesssim 1.8 \times 10^{-8}$ [26] for $m_{DM} \simeq 3.5 \text{ keV}$ [78], one gets
\[ \frac{m_{DM}}{1 \text{ GeV}} \gtrsim 4 \times 10^{-6} \left( \frac{M_{in}}{M_{P}} \right)^{1/2} \simeq 2 \times 10^{-3} \left( \frac{M_{in}}{g} \right)^{1/2}. \] (3.9)

This bound is shown in blue in Fig. 1, and only constrains light DM particles with mass $m_{DM} \lesssim 30 \text{ GeV}$. 

4 Self-interactions

In the previous section, the production of collisionless DM particles via the evaporation of PBHs was presented. However, DM can feature sizable self-interactions, dramatically changing its expected relic density. To analytically understand the role played by DM self-interactions [49], let us study the DM production under the assumption of an instantaneous evaporation of the PBHs at $T = T_{ev}$.

4.1 Radiation Dominated Universe

First, we focus on the case where the universe was dominated by SM radiation during the whole lifetime of PBHs.

4.1.1 Light DM

In the case where the DM is lighter than the initial BH temperature ($m_{DM} \ll T_{BH}^{in}$), each BH radiates $n_{DM}$ DM particles with a mean energy $\langle E \rangle = 6 T_{BH}$, Eq. (2.11). The total DM energy density radiated by a BH can be estimated by
\[ \rho_{DM}(T = T_{ev}) \simeq n_{DM}(T = T_{ev}) T_{BH} = \beta \frac{g_{DM} \zeta(3) C_n}{2 \pi^2} T_{in} T_{ev}^3, \] (4.1)
where the SM entropy conservation and Eq. (2.10) were used. Let us notice that the produced DM population inherits the BH temperature, and, as DM is not in chemical equilibrium, develops a large chemical potential.

If the dark sector has sizable self-interactions guaranteeing that elastic scatterings within the dark sector reach kinetic equilibrium, DM thermalizes with a temperature $T'$ in general different from the SM one $T$ and the BH temperature $T_{BH}$. Assuming an instantaneous thermalization process, and taking into account instantaneous conservation of the DM energy density, the temperature $T'_{ev}$ in the dark sector just after thermalization is

$$T'_{ev} = \begin{cases} \left( \frac{15(3)}{4} C_n \right)^{1/4} T_{in}^{1/4} T_{ev}^{3/4} & \text{for } m_{DM} \ll T'_{ev}, \\ \frac{2}{3} m_{DM} W^{-1} \left[ \frac{1}{\pi^2} \left( \frac{g_{DM} m_{DM}^4}{\rho_{DM}(T=0)} \right)^{2/3} \right] & \text{for } m_{DM} \gg T'_{ev}, \end{cases} \quad (4.2)$$

for the cases $m_{DM} \ll T'_{ev} < T_{BH}^m$ and $T'_{ev} \ll m_{DM} < T_{BH}^m$, respectively, where $W$ corresponds to the principal value of the Lambert function, and $C_n = 1$ (bosonic DM) or $7/8$ (fermionic DM). There is a net decrease of the DM mean kinetic energy $T'_{ev}/T_{BH} \ll 1$, and therefore DM can become non-relativistic due to thermalization effects.

Additionally, if number-changing self-interactions within the dark sector reach chemical equilibrium, the DM number density just after thermalization is therefore

$$n_{DM}(T'_{ev}) = \begin{cases} \frac{15(3)}{4} g_{DM} C_n \left( \frac{1}{\pi^2} \frac{g_{DM} m_{DM}^4}{\rho_{DM}(T=0)} \right)^{3/4} \beta^{3/4} \left( T_{in} T_{ev}^3 \right)^{3/4} & \text{for } m_{DM} \ll T'_{ev}, \\ \frac{2}{3} m_{DM} W^{-1} \left[ \frac{1}{\pi^2} \left( \frac{g_{DM} m_{DM}^4}{\rho_{DM}(T=0)} \right)^{2/3} \right] & \text{for } m_{DM} \gg T'_{ev}. \end{cases} \quad (4.3)$$

The overall effect of self-interactions and in particular of number-changing interactions within the dark sector is to decrease the DM temperature, increasing the DM number density. Such an increase can be characterized by a boost factor $B$ defined by comparing the DM number densities taking into account the case with relative to the case without thermalization in the dark sector:

$$B \equiv \frac{n_{DM}(T'_{ev})}{n_{DM}(T = T_{ev})} \simeq \begin{cases} \frac{15(3)}{4} g_{DM} C_n \left( \frac{1}{\pi^2} \frac{g_{DM} m_{DM}^4}{\rho_{DM}(T=0)} \right)^{3/4} \beta^{1/4} T_{BH}^{1/4} T_{ev}^{3/4} & \text{for } m_{DM} \ll T'_{ev}, \\ \frac{2}{3} m_{DM} W^{-1} \left[ \frac{1}{\pi^2} \left( \frac{g_{DM} m_{DM}^4}{\rho_{DM}(T=0)} \right)^{2/3} \right] & \text{for } m_{DM} \gg T'_{ev}. \end{cases} \quad (4.4)$$

### 4.1.2 Heavy DM

In the case where DM is heavier than the initial BH temperature ($m_{DM} \gg T_{BH}^m$), PBHs radiate DM particles with a mean energy $\langle E \rangle = 6 m_{DM}$, Eq. (2.11). The total DM energy density radiated by a BH can be estimated by

$$\rho_{DM}(T = T_{ev}) \simeq \beta \frac{g_{DM} C_n}{2 \pi^2} \frac{M_{in} m_{DM} m_{DM}^2}{T_{in} T_{ev}} \left( T_{in} T_{ev}^3 \right), \quad (4.5)$$

and therefore, the temperature in the dark sector just after thermalization of non-relativistic DM particles is

$$T'_{ev} \simeq \frac{2}{3} m_{DM} W^{-1} \left[ \frac{1}{\pi^2} \left( \frac{2 \pi^2}{\beta (3)} C_n \frac{m_{DM}^3 M_{in}^2}{T_{in} T_{ev}^3 M_{in}} \right)^{2/3} \right], \quad (4.6)$$
which this time corresponds to a mild decrease of the mean DM kinetic energy $T'_{\text{ev}}/m_{\text{DM}} \lesssim 1$. The DM number density just after thermalization for non-relativistic DM particles is

$$n_{\text{DM}}(T'_{\text{ev}}) = \beta \frac{g_{\text{DM}} \zeta(3) C_n}{2\pi^2} \frac{M_{\text{in}}}{M_P} T_{\text{in}}^3 T_{\text{ev}}^3,$$  \hspace{1cm} (4.7)

and hence a boost factor

$$B \simeq 1.$$  \hspace{1cm} (4.8)

We notice that this factor coincides with the one for DM heavier than $T'_{\text{ev}}$ in Eq. (4.4), by taking the limit $m_{\text{DM}} \to \bar{T}_{\text{BH}}$. A boost factor $B \simeq 1$ (i.e. no boost!) was expected in this case where the originally Hawking radiated particles were almost non-relativistic.

### 4.2 Matter Dominated Universe

After having studied the case where the universe was dominated by SM radiation during the whole lifetime of the PBH, in this section we focus on the other scenario, in which PBHs dominated the energy density. In this section an analysis analogous to the one presented previously will be followed.

#### 4.2.1 Light DM

The total DM energy density for light particles ($m_{\text{DM}} \ll T'_{\text{ev}}$) radiated by a BH is

$$\rho_{\text{DM}}(T = T_{\text{ev}}) \simeq \frac{g_{\text{DM}} \zeta(3) C_n}{2\pi^2} \bar{T}_{\text{ev}}^4.$$  \hspace{1cm} (4.9)

Again, assuming an instantaneous thermalization process, and taking into account the instantaneous conservation of the DM energy density, the temperature $T'_{\text{ev}}$ in the dark sector just after thermalization

$$T'_{\text{ev}} \simeq \begin{cases} \left( \frac{15 \zeta(3) C_n}{\pi^2 \rho_{\text{DM}}(T=\bar{T}_{\text{ev}})} \right)^{1/4} \bar{T}_{\text{ev}} & \text{for } m_{\text{DM}} \ll T'_{\text{ev}}, \\ \frac{2}{\pi m_{\text{DM}}} W^{-1} \left[ \frac{1}{3} \left( \frac{g_{\text{DM}} m_{\text{DM}}^4}{\rho_{\text{DM}}(T=\bar{T}_{\text{ev}})} \right)^{2/3} \right] & \text{for } m_{\text{DM}} \gg T'_{\text{ev}}, \end{cases}$$  \hspace{1cm} (4.10)

the DM number density just after thermalization

$$n_{\text{DM}}(T'_{\text{ev}}) \simeq \begin{cases} \frac{\zeta(3) C_n}{\pi^2} g_{\text{DM}} \left( \frac{15 \zeta(3) C_n}{\pi^2 C_\rho} \right)^{3/4} T_{\text{ev}}^3 & \text{for } m_{\text{DM}} \ll T'_{\text{ev}}, \\ g_{\text{DM}} \zeta(3) C_n \left( \frac{T_{\text{ev}}^3}{m_{\text{DM}}} \right) & \text{for } m_{\text{DM}} \gg T'_{\text{ev}}, \end{cases}$$  \hspace{1cm} (4.11)

and therefore the boost factor becomes

$$B \simeq \begin{cases} \frac{2}{\pi} \left( \frac{15 \zeta(3) C_n}{\pi^2 C_\rho} \right)^{3/4} \bar{T}_{\text{BH}} T_{\text{ev}} & \text{for } m_{\text{DM}} \ll T'_{\text{ev}}, \\ \frac{T_{\text{in}}}{\bar{T}_{\text{BH}}} & \text{for } m_{\text{DM}} \gg T'_{\text{ev}}, \end{cases}$$  \hspace{1cm} (4.12)

which matches the result for the radiation domination case, Eq. (4.4), in the limit $\beta \to T_{\text{ev}}/T_{\text{in}}$ with $\bar{T}_{\text{ev}} = T_{\text{ev}}$. 


4.2.2 Heavy DM

The total DM energy density for heavy particles \(m_{\text{DM}} \gg T_{\text{in}}\) radiated by a BH is

\[
\rho_{\text{DM}}(T = \bar{T}_{\text{ev}}) \simeq \frac{g_{\text{DM}} \zeta(3) C_n}{2\pi^2} \frac{m_{\text{DM}}}{\bar{T}_{\text{BH}}^{4/3}}. \tag{4.13}
\]

The temperature \(T'_{\text{ev}}\) in the dark sector just after thermalization of non-relativistic DM particles

\[
T'_{\text{ev}} \simeq \frac{2}{3} m_{\text{DM}} W^{-1} \left[ \frac{1}{3\pi} \left( \frac{g_{\text{DM}} m_{\text{DM}}^4}{\rho_{\text{DM}}(T = \bar{T}_{\text{ev}})} \right)^{2/3} \right], \tag{4.14}
\]

the DM number density just after thermalization for non-relativistic DM particles becomes

\[
n_{\text{DM}}(T = \bar{T}_{\text{ev}}) \simeq \frac{g_{\text{DM}} \zeta(3) C_n}{2\pi^2} \frac{T_{\text{in}}^4}{\bar{T}_{\text{BH}}^{4/3}}, \tag{4.15}
\]

and hence the boost factor

\[
B \simeq 1, \tag{4.16}
\]

which matches the results for the radiation dominated case, Eq. (4.8), and for light DM in matter domination, Eq. (4.12), in the limit \(m_{\text{DM}} \rightarrow T_{\text{in}}^{3/3}\).

The impact of DM self-interactions is shown in Fig. 2, for different DM masses. Dotted and solid thick lines correspond to the limiting cases without and with a maximal effect from self-interactions, respectively, in the same parameter space used in Fig. 1. Out of the six regimes presented previously, four are visible in the plots and are described in ascending order for \(T_{\text{in}}\).

- The observed DM abundance can be generated in the case where the PBHs dominated the universe energy density (above the red dotted line), only for heavy DM, i.e. \(m_{\text{DM}} > T_{\text{BH}}\). The DM yield is independent from \(\beta\) (Eq. (3.5)) and there is no boost due to self-interactions (Eq. (4.16)), as can be seen in the lower right panel corresponding to \(m_{\text{DM}} = 10\) TeV. We notice that this scenario is typically excluded by the BBN constraint, and only viable for \(m_{\text{DM}} \gtrsim 10^9\) GeV.
- The case of heavy DM, this time in a radiation dominated scenario, appears when \(m_{\text{DM}} \gtrsim 100\) GeV. This case is visible in the lower right panel, where the DM yield is given by Eq. (3.1) without a significant boost, Eq. (4.8). It follows that \(\beta \propto T_{\text{in}}^{-3}\). This scenario is again typically in tension with the BBN observations, and only viable when \(m_{\text{DM}} \gtrsim 10^5\) GeV.
- The third regime corresponds to \(T_{\text{ev}} \ll m_{\text{DM}} \ll T_{\text{BH}}\), where the DM yield in Eq. (3.1) is boosted by a factor \(T_{\text{BH}}^{3/3}/m_{\text{DM}}\), Eq. (4.4). In this case, \(\beta \propto T_{\text{in}}^{-1}\). We notice that values of \(\beta\) smaller than the ones required in this scenario always produce a DM underabundance.
- The last regime happens for light DM, \(m_{\text{DM}} \ll T_{\text{ev}}\) in a radiation-dominated universe. The DM yield in Eq. (3.1) is boosted by a factor \(\propto T_{\text{BH}}^{1/3}/T_{\text{ev}}^{3/3}\), Eq. (4.4). In this case, the DM abundance requires \(\beta \propto T_{\text{in}}^2\).
We would like to emphasize that the current boost factors have to be understood as the maximum increase of the DM number density due to self-interactions. They can be reached if the number-changing DM self-interactions freeze-out while DM is relativistic (or soon after chemical equilibrium is achieved), to avoid a DM depletion due to cannibalization processes, e.g. $3 \to 2$ or $4 \to 2$ annihilations. Furthermore, we notice that in specific models, chemical equilibrium in the dark sector may not be achieved due to the perturbativity limit, and therefore the maximum boost can not be attained.

Additionally to the increase of the DM number density, self-interactions also reduce the...
DM typical momentum. The bound on the DM mass coming from the possible suppression on the structure formation rate due to its free-streaming length is therefore eased. In the case with self-interactions, the momentum of DM particles in a matter-dominated scenario is \( p_{\text{ev}} \simeq T'_{\text{ev}} \), see Eq. (4.10). Therefore, from Eq. (3.7) one has that

\[
p_0 \simeq \left[ \frac{g_{**}(T_{\text{eq}})}{g_{**}(T_{\text{ev}})} \right]^{1/3} \frac{T_{\text{eq}}}{\Omega_m} \frac{15 \zeta(3) C_n}{\pi^4 C_\rho} \frac{1}{\Omega_R}^{1/4},
\]

for light DM with \( m_{\text{DM}} \ll T'_{\text{ev}} \), and assuming that DM momentum scales like \( a^{-1} \). This can be achieved if \( p \) simply redshifts, i.e. if DM kinetically decouples just after thermalization, or if the kinetic equilibrium is broken when DM is still relativistic, and implies a lower limit on the DM mass

\[
m_{\text{DM}} \gtrsim 4 \text{ keV},
\]

following the same procedure as for the case without self-interactions, in section 3. However, if chemical equilibrium is active when DM is non-relativistic, number-changing interactions enforce DM temperature to fall only logarithmically until they freeze-out [29]. These cannibalization processes rise the DM temperature relative to the SM, increasing the DM momentum and therefore strengthening the bound on the DM mass. In that sense, Eq. (4.18) has to be understood as a minimal lower bound.

5 Conclusions

DM production via Hawking evaporation of PBHs constitutes an irreducible process in the early universe. This channel can be dominant for instance, if the dark and visible sectors are disconnected. In that case, after production the DM comoving density stays constant until today. Additionally, light DM is radiated relativistically and could erase small-scale structures via free-streaming. This enforces DM to be heavier than a few MeV.

However, this paradigm is modified if DM features sizable self-interactions. Thermalization and number-changing processes in the dark sector can have strong impacts, in particular enhancing the DM relic abundance by several orders of magnitude. In this paper we have estimated the boost from general arguments such as the conservation of energy and entropy, independently from the underlying particle physics details of the dark sector. Two main consequences can be highlighted: i) As the DM abundance is increased, a smaller initial energy density of PBHs (encoded in the parameter \( \beta \)) is required. ii) Thermalization in the dark sector decreases the mean DM kinetic energy, relaxing the bound from structure formation and hence, allowing for lighter DM in the keV ballpark.

Before concluding, we note that thermalization and number-changing interactions naturally appear in scenarios where DM features sizable self-interactions. Those DM self-interactions could play a role in the solution of the so-called ‘core vs. cusp’ [79–82] and ‘too-big-to-fail’ problems [83–86] arising at small scales. For this to be the case, the required self-scattering cross section over DM mass needs to be of the order of \( 0.1–2 \text{ cm}^2/\text{g} \) at the scale of dwarf galaxies [87, 88], and smaller than \( 1.25 \text{ cm}^2/\text{g} \) at the scale of galaxy clusters [89]. Additionally, possible imprints may arise from the evaporation of PBHs before the BBN onset. Since gravitons are always within the BH evaporation products and do not reach thermal equilibrium, there would be a corresponding stochastic background of gravitational waves with a frequency in the range above 10 THz [90–92], which is high enough to be within the sensitivity region of ongoing and near future gravity wave experiments [24, 59, 60, 93–96].
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