Asymmetric dark matter abundance including non-thermal production

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Abstract

We investigate the relic abundance of asymmetric dark matter where the asymmetric dark matter is non-thermally produced from the decay of heavier particles in addition to the usual thermal production. We discuss the relic density of asymmetric dark matter including the decay of heavy particles in low-temperature scenarios. Here, we still assume that the Universe is radiation-dominated and there is asymmetry before the decay of heavy particles. We obtain an increased abundance of asymmetric dark matter when there is an additional contribution from the decay of heavier particles. Finally, we find the constraints on the asymmetry factor and annihilation cross-section using Planck data.

Keywords: asymmetric dark matter, relic density, non-thermal production

(Some figures may appear in colour only in the online journal)

1. Introduction

The dark matter problem has undoubtedly been one of the most challenging puzzles for cosmologists and particle physicists in recent years. Although we have the most striking evidence for dark matter, what constitutes dark matter is still a mystery to us. Asymmetric dark matter is one alternative, which arises from the fact that its abundance is just a few times larger than the baryon asymmetry [1]. There may be an indication of a common origin, which is responsible for the baryonic asymmetry and dark matter. In an asymmetric dark matter scheme, dark matter particles have distinct antiparticles.

In a standard cosmological scenario, it is supposed that asymmetric dark matter is produced thermally. Based on the assumption that weakly interacting massive particles constitute dark matter, the relic density of asymmetric dark matter is determined by the freeze-out condition [2, 3]. In this scenario, it is assumed that the reheating temperature \( T_R \) is much larger than the freeze-out temperature \( T_F \) and the asymmetric dark matter particles and antiparticles were in thermal equilibrium in the early Universe. The final relic density of asymmetric dark matter mainly depends on the particle–antiparticle asymmetry, and the indirect detection signals from annihilation are suppressed because there are very few antiparticles here in the end. On the other hand, because we have observational evidence from Big Bang Nucleosynthesis for temperatures \( T \lesssim \text{few} \) MeV [4–6], it is also interesting to explore the asymmetric dark matter in scenarios with \( T_R \lesssim T_F \) [7–9]. Asymmetric dark matter can be produced non-thermally in low reheating temperature scenarios since heavy particles decay (e.g. moduli decay) into asymmetric dark matter. In [7], the authors discuss a cogenesis mechanism in which the baryon asymmetry of the Universe and dark matter abundance are simultaneously produced at low reheating temperature. Reference [8] explored the non-thermal production of dark matter including asymmetric dark matter in the scheme of the Minimal Supersymmetric Standard Model and string-motivated models. Relic abundance of dark matter was calculated in [10] including the decay of unstable heavy particles to dark matter particles for low reheating temperature.

In our work, we investigate the relic density of asymmetric dark matter in a low reheating temperature scenario when the heavy unstable particles \( \phi \) decay into asymmetric dark matter particles and antiparticles. We suppose that the asymmetry already existed before the decay of heavy particles into asymmetric dark matter and that they decay into particles and antiparticles in the same amount. We assume that the metastable particles never dominate the energy density of the Universe. There are several issues that are not present in the radiation-dominated epoch in low reheating temperature scenarios since the entropy density is not conserved and non-perturbative inflaton decays may be significant [11–13]. We avoid these problems by considering the abundance of asymmetric dark...
matter at some initial temperature $T_0$ as a free parameter. We find that the decay of heavier particles to asymmetric dark matter changes the total relic density of asymmetric dark matter. The relic abundances of asymmetric dark matter particles and antiparticles are both increased when there is an additional contribution from the decay of heavier particles.

The outline of the paper is as follows. In section 2, we discuss the relic density calculation of asymmetric dark matter including the decay of heavy particles into asymmetric dark matter in low-temperature scenarios. Here, we assume that the Universe is still radiation-dominated and additional entropy production from the decay of heavy particles is negligible. We find the approximate analytical solution for the relic density of asymmetric dark matter when there is thermal and non-thermal production. In section 3, the constraints on parameter spaces are obtained by using Planck data. The last section is devoted to a brief summary and conclusions.

2. Relic abundance of asymmetric dark matter including non-thermal production

We discuss the scenario where the unstable heavy particles $\phi$ decay into asymmetric dark matter particles and antiparticles. Here, we assume that the heavy-particle $\phi$ decays out of thermal equilibrium. Therefore, $\phi$ production is negligible. However, we include the thermal and non-thermal production of asymmetric dark matter particles.

In the standard computation of relic density for the asymmetric dark matter, it was assumed that the reheating temperature of the Universe $T_R$ is much higher than the freeze-out temperature. In this case, the reheating era has no effect on the final relic density of asymmetric dark matter. On the other hand, the constraints on the reheating temperature originate from Big Bang Nucleosynthesis and $T_R \gtrsim 1$ MeV [4–6]. Therefore, we consider the case where $T_R < T_f$ [7–9]. The asymmetric dark matter particles and antiparticles never reach thermal equilibrium due to the low reheating temperature. Thus, we treat the asymmetric dark matter abundance at some initial temperature $T_0$ as a free parameter.

The evolution of relic abundance including the decay of heavy particles is more complicated than the usual thermal production case. Here, we still assume that the Universe is radiation dominated and the entropy production is negligible. We suppose that asymmetry existed before the decay of heavy particles and they decay into particles and antiparticles in the same amount. The coupled Boltzmann equations for the number densities $n_\chi$ and $n_{\tilde{\chi}}$ of asymmetric dark matter particles $\chi$ and antiparticles $\tilde{\chi}$ are as follows:

$$\frac{dn_\chi}{dt} + 3Hn_\chi = -\langle \sigma v \rangle (n_\chi n_{\tilde{\chi}} - n_{\chi,eq}n_{\tilde{\chi},eq}) + N\Gamma_0 n_\phi,$$  
(1)

$$\frac{dn_{\tilde{\chi}}}{dt} + 3Hn_{\tilde{\chi}} = -\langle \sigma v \rangle (n_\chi n_{\tilde{\chi}} - n_{\chi,eq}n_{\tilde{\chi},eq}) + N\Gamma_0 n_\phi,$$  
(2)

$$\frac{dn_\phi}{dt} + 3Hn_\phi = -\Gamma_0 n_\phi,$$  
(3)

where $H$ is the expansion rate of the Universe, and $\langle \sigma v \rangle$ is the thermal average of the cross-section multiplied by the relative velocity of annihilating asymmetric dark matter particles and antiparticles. Here, we assume that only $\chi\tilde{\chi}$ pairs can be annihilated into Standard Model particles. $n_{\chi,eq}$ is the equilibrium number density, $N$ is the average number of asymmetric dark matter particles and antiparticles produced in $\phi$ decay, while $\Gamma_0$ is the decay rate. Here, it is assumed that $\phi$ does not dominate the total energy density. Therefore, the comoving entropy density stays almost constant throughout. The analytical solution for equation (3) is easily obtained. Then, equations (1) and (2) can be solved approximately in low-temperature scenarios.

In order to simplify equations (1), (2) and (3), we introduce the following dimensionless variables $Y_\chi, 4 = n_{\chi,eq}/s, Y_0 = n_\phi/s$ and $x = m/T$ with $s$ being the entropy density, then:

$$\frac{d(sY_\chi)}{dt} + 3sY_\chi = -\langle \sigma v \rangle s^2(Y_\chi Y_\tilde{\chi} - Y_{\chi,eq}Y_{\tilde{\chi},eq}) + N\Gamma_0 sY_0,$$  
(4)

$$\frac{d(sY_{\tilde{\chi}})}{dt} + 3sY_{\tilde{\chi}} = -\langle \sigma v \rangle s^2(Y_\chi Y_\tilde{\chi} - Y_{\chi,eq}Y_{\tilde{\chi},eq}) + N\Gamma_0 sY_0,$$  
(5)

$$\frac{d(sY_\phi)}{dt} + 3sY_\phi = -\Gamma_0 sY_\phi.$$  
(6)

Assuming that entropy per comoving volume is conserved, we obtain:

$$\dot{s} = -3sH,$$  
(7)

where we used $H = \dot{R}/R$, where $R$ is the scale factor of the Universe. Inserting $s = (2\pi^2/45)g_*T^3$ into equation (7), where $g_*$ is the effective number of relativistic degrees of freedom, we obtain:

$$\frac{dT}{dt} + \frac{1}{3g_*} \frac{dg_*}{dt} + H = 0,$$  
(8)

Using $T = m/x$:

$$\frac{dx}{dt} = -\frac{Hx}{\frac{1}{3} - \frac{1}{3g_*}\frac{dg_*}{dx}},$$  
(9)

Then, equations (4), (5) and (6) can be written as,

$$\frac{dY_\chi}{dx} = -\frac{\langle \sigma v \rangle s}{Hx}(Y_\chi Y_\tilde{\chi} - Y_{\chi,eq}Y_{\tilde{\chi},eq}) + N\Gamma_0 Y_0,$$  
(10)

$$\frac{dY_{\tilde{\chi}}}{dx} = -\frac{\langle \sigma v \rangle s}{Hx}(Y_\chi Y_\tilde{\chi} - Y_{\chi,eq}Y_{\tilde{\chi},eq}) + N\Gamma_0 Y_0,$$  
(11)

$$\frac{dY_\phi}{dx} = -\frac{\Gamma_0}{Hx} Y_\phi,$$  
(12)

where we assume that $dg_*/dx = 0$. Equation (12) can then be easily solved:

$$\int_{Y_{\phi}(x_0)}^{Y_{\phi}(x)} \frac{dY_\phi}{Y_\phi} = -\int_{x_0}^{x} \frac{\Gamma_0 dx}{Hx},$$  
(13)

$$Y_0(x) = Y_0(x_0)e^{-\int_{x_0}^{x} \frac{\Gamma_0 dx}{Hx}},$$  
(14)

where $x_0$ is the inverse scaled initial temperature, which is close to the freeze-out temperature. Then, equations (10) and (11)
In the radiation-dominated era, $H = \pi m_\chi^2/(M_{Pl} x^2)^{1/2}$, where $m_\chi$ is the mass of asymmetric dark matter particles and $M_{Pl} = 2.4 \times 10^{18}$ GeV is the reduced Planck mass. We then further simplify equations (15) and (16), subtracting equation (16) from (15), then,

$$\frac{d(Y_\chi - Y_\bar{\chi})}{dx} = 0. \tag{17}$$

This requires the following:

$$Y_\chi - Y_\bar{\chi} = \eta, \tag{18}$$

where $\eta$ is a constant. We can express equations (15) and (16) in terms of equation (18):

$$\frac{dY_\chi}{dx} = -\frac{\langle \sigma v \rangle}{H x^2} (Y_\chi - \eta Y_\bar{\chi} - \lambda_2 x^2 e^{-2x})$$

$$+ \lambda_3 N Y_\bar{\chi}(x_0) \exp \left[ \frac{\lambda_2}{2} (x_0^2 - x^2) \right], \tag{19}$$

$$\frac{dY_\bar{\chi}}{dx} = -\frac{\langle \sigma v \rangle}{H x^2} (Y_\bar{\chi} - \eta Y_\chi - \lambda_2 x^2 e^{-2x})$$

$$+ \lambda_3 N Y_\chi(x_0) \exp \left[ \frac{\lambda_2}{2} (x_0^2 - x^2) \right], \tag{20}$$

where $\lambda_2 = \lambda_1^2/(\pi m_\chi^2/m_{Pl}^2)^{1/2}$ and $\lambda_3 = M_{Pl}/(\pi m_\chi^2)^{1/2}$. Here, we used the equilibrium abundance $Y_{\chi,\bar{\chi}, eq} = 0.145 g_\chi^2/x^3 e^{-x(1+\mu_\chi/m_\chi)}$, where $\mu_\chi$ is the chemical potential for particles, and $g_\chi$ is the number of internal degrees of freedom. The thermal average of $\sigma v$ is usually approximated by a nonrelativistic expansion:

$$\langle \sigma v \rangle = a + 6 bx^{-1} + O(x^{-2}), \tag{21}$$

where $a$ is the $s$-wave contribution for the limit $v \to 0$ and $b$ is the $p$-wave contribution for the suppressed $s$-wave annihilation.

Figure 1 depicts the evolution of relic abundance $Y_\chi,\bar{\chi}$ as a function of $x$ and it is based on the numerical solutions of equations (19) and (20). Here, the solid red curve is the asymmetric dark matter particle abundance $Y_\chi, nt$ with the decay of...
Figure 2. Relic density $\Omega h^2$ for particle $\chi$ and antiparticle $\bar{\chi}$ as a function of the $s$-wave cross-section. Here, $Y_i(x_0 = 22) = 1 \times 10^{-12}, N = 1$ and $Y_o = 10^{-11}, \Gamma = 10^{-11}, \eta = 3 \times 10^{-12}, m_\chi = 100 \text{GeV}$. $x_0 = 22, g_y = 2, g_a = 90$.

heavy particles, and the black dotted curve is the antiparticle abundance $Y_{\bar{\chi}, \text{er}}$. The red dot-dashed and black dashed curves are for asymmetric dark matter particle and antiparticle relic abundances $Y_i$ and $Y_{\bar{\chi}}$ without including the decay of heavy particles. Figure 1 shows that the abundances of asymmetric dark matter particles and antiparticles are increased due to the contribution from the decay of metastable heavier particles to asymmetric dark matter. The increase in antiparticle abundance is more remarkable. We start from the initial value of $Y_{\bar{\chi}}(x_0)$ at the initial temperature $x_0$. We assume that asymmetric dark matter particles and antiparticles never reach thermal equilibrium due to the low reheating temperature. The particle and antiparticle abundances are increased quickly, shortly after starting, and then decreased both for thermal and non-thermal production. When $x$ increases, the abundance finally goes to a constant value. The increase in particle and antiparticle abundances is slightly larger in panels (b) and (d) of figure 1 for a smaller decay rate with respect to panels (a) and (c). Comparing frames (a) with (c), and (b) with (d), it can be seen that for the larger initial value of $Y_{\chi}(x_0)$, there is a larger increase in the particle and antiparticle abundances. In frames (c) and (d), the antiparticle abundance almost reaches the same level as that of the particle abundance. This opens the possibility to detect the asymmetric dark matter indirectly.

The changes in relic density $\Omega h^2$ as a function of the cross-section are shown in figure 2. Here, $s$-wave annihilation cross-section with $b = 0$ when asymmetry factor $\eta = 3 \times 10^{-12}$. The solid red curve and dotted black curve represent the relic densities of asymmetric dark matter particles and antiparticles with the decay of heavy particles. The double-dotted red curve and black dashed curve represent the particles and antiparticles only with thermal production. Figure 2 shows when the cross-section increases. Here, the depletion of antiparticle abundance is slower than the case without including heavy-particle decay.

In what follows, we try to find the approximate analytical solutions of equations (19) and (20). Equation (20) can be solved first. Initially, $(x < x_0), Y_{\chi, \text{eq}} \gg Y_{\chi}$, because $\eta$ is usually around $10^{-12}$ [3], and $\eta Y_{\chi}$ is small as well. Therefore, equation (20) becomes,

$$\frac{dY_{\chi}}{dx} = \lambda_1 \lambda_2 (\sigma v) x e^{-2x} + \lambda_1 Y_o(x_0)x \exp \left( \frac{\lambda_1}{2} (x_0^2 - x^2) \right).$$

Integrating the above equation, we obtain,

$$Y_{\chi}(x) = Y_i(x_0) + \lambda_1 \lambda_2 \left[ \frac{a}{2} (x_0 - x) \exp \left( \frac{a}{4} (x_0^2 - x^2) \right) + \left( \frac{a}{4} + 3b \right) (x_0^2 - x^2) \right] + NY_o(x_0) \left[ 1 - \exp \left( \frac{\lambda_1}{2} (x_0^2 - x^2) \right) \right].$$

The abundance for $Y_{\bar{\chi}}$ is obtained using equation (18):

$$Y_{\bar{\chi}}(x) = \eta + Y_{\bar{\chi}}(x_0) + \lambda_1 \lambda_2 \left[ \frac{a}{2} (x_0 - x) \exp \left( \frac{a}{4} (x_0^2 - x^2) \right) + \left( \frac{a}{4} + 3b \right) (x_0^2 - x^2) \right] + NY_o(x_0) \left[ 1 - \exp \left( \frac{\lambda_1}{2} (x_0^2 - x^2) \right) \right].$$

While the solution is valid only initially, the relic abundances of particles and antiparticles are approximated for $x \gg x_0$ as,

$$\begin{align*}
Y_{\chi, \infty} &\equiv Y_i(x \gg x_0) = Y_i(x_0) + \lambda_1 \lambda_2 \times \left[ \frac{a}{2} x_0 \exp \left( \frac{a}{4} (x_0^2 - x_0^2) \right) + \left( \frac{a}{4} + 3b \right) (x_0^2 - x_0^2) \right] + NY_o(x_0),
\end{align*}$$

$$\begin{align*}
Y_{\bar{\chi}, \infty} &\equiv Y_{\bar{\chi}}(x \gg x_0) = \eta + Y_{\bar{\chi}}(x_0) + \lambda_1 \lambda_2 \times \left[ \frac{a}{2} x_0 \exp \left( \frac{a}{4} (x_0^2 - x_0^2) \right) + \left( \frac{a}{4} + 3b \right) (x_0^2 - x_0^2) \right] + NY_o(x_0).
\end{align*}$$

The relative prediction for the present relic density is given by,

$$\Omega_{\text{DM}} h^2 = \frac{\rho}{\rho_c} = 0.9 \times 10^6 m_\chi \left[ Y_{\chi, \infty} + Y_{\bar{\chi}, \infty} \right] \text{GeV}^{-1},$$

where the scaled Hubble constant $h$ in units of 100 km s$^{-1}$ Mpc$^{-1}$ is 0.673 ± 0.098. Here, $\rho = \text{nm} = s_0 Y m$ and the critical density is $\rho_c = 3H_0^2 M_\odot h^2$, where $s_0 \approx 2900$ cm$^{-3}$ is the present entropy density and $H_0$ is the Hubble constant.

In the late Universe, for $x \gg x_0$,

$$Y_{\chi} \gg Y_{\chi, \text{eq}}.$$ (28)

Again, the third and fourth terms in equation (20) decrease exponentially as $x$ increases. Therefore, equation (20) becomes,

$$\frac{dY_{\chi}}{dx} = -\frac{\lambda_1 (\sigma v)}{x^2} (Y_{\chi}^2 + \eta Y_{\chi}).$$

Integrating the above equation, we obtain,

$$Y_{\chi}(x) = \frac{\eta}{[1 + \eta Y_{\chi}(x_0)] e^{\int_0^x \frac{\lambda_1 (\sigma v)}{x^2} dx}}.$$ (30)
same as in $\Omega$ tions of the Boltzmann equations (We use the measured dark matter relic density to
The Planck team derived the dark matter relic density as $\Omega_{\text{DM}}h^2 = 0.1199$. Other parameters are the same as in figure 2.

In the same way as equation (24), we have,

$$Y_\chi(x) = \frac{\eta}{1 - \{Y_\chi(x_0)/[\eta + Y_\chi(x_0)]\}e^{-\int_{x_0}^{x} \frac{\Omega_{\text{DM}}m_{\chi}}{T^2}}.}$$

(31)

Inserting equation (21) into (30) and (31), we obtain,

$$Y_\chi(x) = \frac{[1 + \eta/Y_\chi(x_0)] \exp \{\eta \lambda_1[a(1/x_0 - 1/x) + 3b(1/x_0^2 - 1/x^2)]\} - 1}{1 - \{Y_\chi(x)/[\eta + Y_\chi(x_0)]\} \exp \{\eta \lambda_1[a(1/x_0 - 1/x) + 3b(1/x_0^2 - 1/x^2)]\}}.$$

$$Y_\bar{\chi} = \frac{\eta}{1 - \{Y_\bar{\chi}(x_0)/[\eta + Y_\bar{\chi}(x_0)]\} \exp \{\eta \lambda_1[a(1/x_0 - 1/x) + 3b(1/x_0^2 - 1/x^2)]\}}.$$

The final relic density is as follows:

$$\Omega_{\text{DM}}h^2 \simeq 2.76 \times 10^8 m_\chi[Y_\chi(x) + Y_\bar{\chi}(x)].$$

3. Constraints on parameter space

The Planck team derived the dark matter relic density as [14],

$$\Omega_{\text{DM}}h^2 = 0.1199 \pm 0.0022.$$  

(35)

We use the measured dark matter relic density to find constraints on the parameter space. For asymmetric dark matter, particle and antiparticle contributions have to be added as $\Omega_{\text{DM}} = \Omega_\chi + \Omega_{\bar{\chi}}$.

The contour plot of the $s$-wave annihilation cross-section $a$ and asymmetry factor $\eta$ when $\Omega_{\text{DM}}h^2 = 0.1199$. This figure is based on the numerical solutions of the Boltzmann equations (19) and (20). The dotted (black) and dashed (black) curves represent the case where the decay of heavier particles to asymmetric dark matter is included. The solid (red) curve represents the standard case. We find that the required cross-section for $Y_\chi(x_0) = 10^{-11}$ is nearly one order of magnitude larger than the standard one in order to fall into the observational range. The increased relic density of asymmetric dark matter due to the decay of heavier particles is suppressed by the larger cross-section. Furthermore, the figure shows that the required cross-section is slightly larger when the initial value of $Y_\chi$ is large. We note that the horizontal part of the contours is not affected by the decay of heavier particles. The antiparticle dark matter density of $Y_\bar{\chi}$ is exponentially suppressed when the asymmetry increases and the total dark matter relic density is determined by the particle density as $\Omega_{\text{DM}} \sim m_{Y_\chi} \sim m_{\chi}$. Therefore, the final relic density of dark matter is independent of cross-section $a$.

4. Summary and conclusion

The relic abundance of asymmetric dark matter, including the decay of heavier particles, is discussed in low reheating temperature scenarios. Here, we assume that there is asymmetry before the unstable heavy particles decay into asymmetric dark matter particles and antiparticles in the same amount. We found that relic abundances of asymmetric dark matter particles and antiparticles are both increased when there is an additional contribution from the decay of heavier particles. The increase in particle and antiparticle abundances due to the non-thermal production depends on the decay rate of heavy particles and slightly on the initial value of heavy-particle abundance. Comparison with the usual thermal production shows that the depletion of the antiparticle abundance is slower for the case of including decay of heavy particles.

Finally, we used the observed dark matter abundance to obtain the constraints on the annihilation cross-section and asymmetry factor when there is a contribution from the decay of heavier particles to asymmetric dark matter. We found that the required annihilation cross-section with the decay of heavy particles is almost one order of magnitude larger than the case without including non-thermal production. These results are important for our understanding of asymmetric dark matter. It is possible to detect asymmetric dark matter through indirect detection due to the increased amount of antiparticle relic density when there is an additional contribution from the decay of heavy particles.

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