Sommerfeld enhancement of invisible dark matter annihilation in galaxies and galaxy clusters

Man Ho Chan

Department of Science and Environmental Studies, The Hong Kong Institute of Education

Abstract

Recent observations indicate that core-like dark matter structures exist in many galaxies, while numerical simulations reveal a singular dark matter density profile at the center. In this article, I show that if the annihilation of dark matter particles gives invisible sterile neutrinos, the Sommerfeld enhancement of the annihilation cross-section can give a sufficiently large annihilation rate to solve the core-cusp problem. The resultant core density, core radius, and their scaling relation generally agree with recent empirical fits from observations. Also, this model predicts that the resultant core-like structures in dwarf galaxies can be easily observed, but not for large normal galaxies and galaxy clusters.

Keywords: Dark Matter

1. Introduction

It is commonly believed that the existence of dark matter can account for the missing mass in galaxies, galaxy clusters and our universe. However, the nature of dark matter remains a fundamental problem in astrophysics. If the dark matter particles are cold and collisionless, $N$-body simulations show that the density profile should be singular at the center (a cusp profile, $\rho \sim r^{-1}$) [1]. This model generally gives good agreements with observations on large-scale structures such as Ly$\alpha$ spectrum [2, 3] and some galaxy clusters [4]. However, observations reveal that cores exist in many galaxies ($\rho \sim r^{-\gamma}$ with $\gamma < 0.5$), especially in dwarf galaxies [5, 6, 7]. Some dwarf galaxies can even have $\gamma < 0.2$ [8]. This discrepancy is commonly known as the core-cusp problem [3]. Moreover, recent studies show that this problem might also be
associated with another problem, called the too-big-to-fail (TBTF) problem. This problem illustrates the fact that the densities of dark matter subhaloes which surround nearby dwarf spheroidal galaxies are significantly lower than those of the most massive subhaloes expected around a normal sized galaxies in cosmological simulations [9, 10]. In other words, solving the core-cusp problem might also provide a solution to the TBTF problem [11].

Some proposals have been suggested to solve the core-cusp problem. For example, the existence of keV dark matter particles, as a candidate of warm dark matter (WDM), has been proposed to solve the problem [12, 13]. However, recent observations indicate that the simplest model of WDM (e.g. the non-resonant sterile neutrino model) cannot account for the major component of dark matter [14, 15, 16, 17]. Some extra properties of WDM or free parameters are needed in order to satisfy the observational constraints. Some recent analyses even suggest that WDM model cannot solve the core-cusp problem [18]. Another proposal suggests that core-like structures would be produced if dark matter particles are self-interacting [19]. Simulations show that dark matter particles with a constant cross-section per unit mass \( \sigma/m \sim 1 \text{ cm}^2 \text{ g}^{-1} \) can produce core-like structures in galaxies [20, 21]. Unfortunately, recent observations put a tight constraint on the cross-section: \( \sigma/m \leq 1 \text{ cm}^2 \text{ g}^{-1} \) [22, 23]. Therefore, it leaves only a small window open for this velocity-independent self-interacting dark matter model to work [23]. Nevertheless, this model might still endure if the cross-section is velocity-dependent, though some more parameters have to be involved [24, 25].

Besides the above two proposals, some suggest that the energy exchange between baryons and dark matter particles might also be possible to produce core-like structures. These mechanisms include the steller and supernova feedback [26, 27], and dynamical friction [28, 29]. It is now a controversial issue because these baryonic processes involve some uncertainties, such as the total energy released by the supernovae and the fraction of energy that can be transferred to the dark matter haloes [11]. For example, for a total mass of \( 10^9 M_\odot \), recent studies show that at least 1/20 of supernova energy is required to transfer to the dark matter halo to give \( \gamma < 0.6 \) [11, 30]. However, we are not sure whether this fraction of transferred energy is physically possible or not. Also, it is challenging to invoke baryonic processes as the main mechanisms to solve the core-cusp problem for some dark-matter-dominated galaxies because the baryonic content is too small to affect the dark matter distribution [8, 25].

In this article, I suggest another proposal that the core-cusp problem
can be solved by Sommerfeld-enhanced dark matter annihilation. The possibility of the dark matter annihilation to solve the core-cusp problem is first suggested in [31]. The required annihilation cross-section is $< \sigma v > \sim 10^{-19} (m/\text{GeV}) \text{ cm}^3 \text{ s}^{-1}$ [31]. They suggest two possible mechanisms so that the required cross-section would not violate the relic dark matter annihilation cross-section $< \sigma v > = 3 \times 10^{-26} \text{ cm}^3 \text{ s}^{-1}$ [31]. However, if the annihilation products are visible particles such as neutrinos, electrons or photons, this required cross-section is ruled out by observations [32, 33]. Moreover, this model predicts that halo core density is universal ($\sim 0.02 M_\odot \text{ pc}^{-3}$) [31] while observations indicate that the core density of dwarf galaxies varies from 0.01 – 0.1 $M_\odot \text{ pc}^{-3}$ [34]. In the following, I use the idea from [31] but assume that the dark matter annihilation is enhanced by the Sommerfeld’s mechanism ($< \sigma v > \propto v^{-\alpha}$) [35, 36], and the annihilation products are invisible sterile neutrinos only. It can be shown that a significant amount of dark matter particles would be annihilated, which is enough to produce the observed core-like structures in galaxies.

2. The annihilation model

It has been suggested that dark matter would self-annihilate to give smaller particles with high energy. The possible stable products formed are photons, electron-positron pairs, and neutrinos. In particular, the fact that active neutrino have non-zero rest mass probably suggests that right-handed neutrinos should exist, which may indeed be sterile neutrinos [37]. Some recent models suggest that dark matter particles can annihilate dominantly into light dark neutrinos (sterile neutrinos) via exchange of a Higgs field ($\chi \chi \to \Phi \Phi \to \nu_s \nu_s$) [38]. This model can agree with the results obtained from DAMA [39] and CoGeNT [40] experiments, which point toward light dark matter ($m \sim 1 – 10 \text{ GeV}$) with isospin-violating and possibly inelastic couplings [38]. However, the light dark matter model is largely constrained by observations, such as cosmic microwave background and SuperK limits. Cline and Frey (2012) [38] propose a model of quasi-Dirac dark matter, interacting via two gauge bosons, one of which couples to baryon number and the other which kinetically mixes with the photon. The annihilation product is dark neutrinos that do not mix with the Standard Model. They also show that the dark neutrinos produced in the universe would not violate the current observational bounds [38].

In the following, we are going to discuss the consequences of the sterile
neutrinos being the only dark matter annihilation product (the model proposed in [38] may be one of the possible scenarios). If the annihilation cross-section is large enough due to the Sommerfeld enhancement, a significant amount of dark matter would be changed to high energy sterile neutrinos. These high energy sterile neutrinos would finally leave the structure and make the central density lower. The Sommerfeld enhancement arises when a scattering object is coupled to a light mediator particle [41]. This enhancement can increase the cross-section for annihilation process in a velocity-dependent fashion due to the generic attractive force between the incident dark matter particles [36].

If dark matter particles were produced at the very beginning, the annihilation cross-section should be close to $3 \times 10^{-26} \text{ cm}^3 \text{ s}^{-1}$ [42]. However, the Sommerfeld enhancement might significantly change the relic annihilation cross-section to a lower value [23, 43, 45]. The actual value of the annihilation cross-section at the thermal freeze out is model-dependent. Nevertheless, [23] show that the Sommerfeld enhancement near resonance would suppress the dark matter abundance by a factor of a few. As a result, the annihilation cross section needs to be suppressed by this same factor in order to be consistent with the observed relic density.

Therefore, we write the annihilation cross section at the thermal freeze out as $\langle \sigma v_0 \rangle = 1 \times 10^{-26} f \text{ cm}^3 \text{ s}^{-1}$, where $f \sim 1$ is a model-dependent parameter. Here, $v_0 \approx 0.2 c$ is the velocity of the dark matter particles at decoupling [46].

If the velocity of dark matter particles $v$ is lower, the Sommerfeld enhancement might increase the cross-section to $\langle \sigma v \rangle = \langle \sigma v_0 \rangle \left(\frac{v_0^\alpha}{v^\alpha}\right)$ [35, 36]. The value of $\alpha$ is close to 1 for non-resonance case while $\alpha \approx 2$ for resonance [44]. Therefore, the low velocity of dark matter particles in a galaxy or galaxy cluster would increase the rate of annihilation to form core-like structures.

Also, this cross-section satisfies the unitarity bound. Harling and Petraki (2014) [47] obtain the unitarity bound (upper bound) of the Sommerfeld-enhanced cross-section, which is just a factor of 3 greater than the unitarity bound without Sommerfeld enhancement. For $m \sim 1 \text{ GeV}$ and $v \sim 10 - 1000 \text{ km/s}$, the unitarity bound is $\langle \sigma v \rangle_{ub} \sim 10^{-13} - 10^{-11} \text{ cm}^3 \text{ s}^{-1}$ [47, 48, 49], which is much greater than the cross-section considered in our model.

On the other hand, the presence of sterile neutrinos as a by-product of dark matter annihilation would change the effective number of neutrinos ($N_{eff}$) in cosmology. This number is strongly constrained by cosmic mi-
crowave background anisotropies recently, which gives $N_{\text{eff}} = 2.88 \pm 0.20$ \cite{50}. However, $N_{\text{eff}}$ depends on the kinetic decoupling temperature of dark matter $T_{\text{de}}$, which is a model-dependent parameter in cosmology. Assuming a reliable range $T_{\text{de}} \sim 10 - 1000 \text{ MeV}$ (at $10^{-7} - 10^{-3} \text{ s}$ after Big Bang) \cite{51}, the number density of active neutrinos is $n_{\nu} \sim 10^{34} - 10^{40} \text{ cm}^{-3}$. The number density of sterile neutrinos produced from dark matter annihilation is given by $n_s \sim \rho_{\text{DM}}^2 < \sigma v_0 > t_{\text{de}}/m^2 \sim 10^{23} - 10^{32} \text{ cm}^{-3}$, where $\rho_{\text{DM}}$ and $t_{\text{de}}$ are the mass density of dark matter and the age of universe at the kinetic decoupling respectively. Therefore, the sterile neutrinos produced at the kinetic decoupling are negligible compared with the relic active neutrinos. Since the scale factor dependence are the same for $n_s$ and $n_{\nu}$, the sterile neutrinos produced would not significantly affect $N_{\text{eff}}$.

Since the mass of the mediator particle $m_\phi$ must be smaller than the dark matter mass \cite{41, 52}, this requires $m_\phi \sim 1 \text{ MeV} - 1 \text{ GeV}$. This new scalar field is likely to mix with the Higgs boson and possibly be in tension with current collider constraints. In view of this, a study in \cite{53} examines this mixing effect seriously by assuming $m_\phi \sim 1 \text{GeV}$ for the Sommerfeld enhancement. Since the mixing between the scalar field and the Higgs field involves some unknown free parameters, including $m_\phi$ and the mixing angle $\theta_{\phi h}$, a large area of parameter space is still possible to satisfy the constraints obtained in the Large Hadron Collider (LHC) experiments \cite{53}. For the CMS measurements, there is an experimental upper limit for the Higgs total width of cross coupling. Based on the calculations in \cite{53}, the upper limit of the Higgs-\phi cross coupling $\lambda_1$ converges to about 0.05 for $m_\phi < 1 \text{ GeV}$, which is not a prohibited value.

Besides the Higgs constraints, the mediator particle would probably emit photons or charged bosons which would possibly conflict with the constraints from indirect detection experiments. However, the emission of photons by the mediator is model-dependent. For example, a study in \cite{54} estimates a corresponding annihilation cross section by considering two dark matter particles annihilating via an intermediate pseudoscalar $A^0$ and a charged fermion $f$ in the loop (see the Feynmann diagram in \cite{54}). For a small coupling ($g \approx 1$), the model gives $< \sigma \gamma \gamma v > \sim 10^{-40} \text{ cm}^3 \text{ s}^{-1} \left( m/1 \text{ GeV} \right)^4 \left( 500 \text{ GeV}/m_A \right)^4 \left( 500 \text{ GeV}/m_f \right)^2$, where $m_A$ and $m_f$ are the mass of the pseudoscalar particle and the fermion respectively \cite{54}. For $m < 1 \text{ GeV}$, this cross-section is well-below any current constraint.
3. Invisible dark matter annihilation in galaxies and galaxy clusters

In fact, numerical simulations should be performed in order to obtain a precise picture of dark matter density profile \( \rho(r, t) \) with annihilation. However, as we will see later, the annihilation is important only when the time \( t \) is sufficiently large (\( t > 1 \text{ Gyr} \)). Therefore, we can simply ignore the dynamics of the halo formation and assume that the initial dark matter density profile follows the NFW profile \([1]\):

\[
\rho(r, 0) = \frac{\rho_s r_s^3}{r(r + r_s)^2},
\]

(1)

where \( \rho_s \) and \( r_s \) are scale density and scale radius of a structure respectively. Numerical simulations show that \( \rho_s r_s \approx 144(M/10^{12} M_{\odot})^{0.2} M_{\odot} \text{pc}^{-2} \) \([55]\), where \( M \) is the total mass of a structure. The time evolution of the dark matter density profile \( \rho(r, t) \) is given by

\[
\frac{\partial \rho(r, t)}{\partial t} = -\left[ \rho(r, t) \right]^2 < \sigma v_0 > v_{\alpha}^0 \frac{m[v(r, t)]^\alpha}{m[v(r, t)]^\alpha}.
\]

(2)

Moreover, in a stable configuration, the circular velocity of dark matter should depend on the mass profile \( M(r, t) \) of a structure: \( v(r, t) \approx \sqrt{GM(r, t)/r} \). According to virial theorem and observations, the velocity dispersion is \( v(r, t) \approx v(r, t) \) \([56]\). Therefore, we assume that the relative velocity of the annihilating dark matter particles follows the velocity profile \( v(r, t) \).

It is believed that most galaxies, including dwarf galaxies, have a supermassive black hole with mass \( M_{BH} \) at the center \([57, 58, 59]\). Therefore we write \( M(r, t) = M_{BH} + M_{DM} \), where \( M_{DM} = \int 4\pi \rho(r, t) r^2 dr \) is the mass profile of dark matter. Here, we do not include the baryonic component because we find that different functional forms of baryonic mass profile just affect the resultant density profile slightly (for example, the difference between using an isothermal and a constant baryonic profile is less than 1% in the resultant density profile). While in galaxy clusters, we have \( M(r, t) = M_{DM} + M_b \), where \( M_b \) is the mass profile of baryonic matter. The baryonic component is important when \( r \) is very small because most of the dark matter is annihilated and there is no supermassive black hole at the centre of galaxy cluster (except cD clusters). In general, the locations of the galactic supermassive black holes in a galaxy cluster are usually far away from the gravitational centre of galaxy cluster (except cD clusters). Therefore, we set \( M_{BH} = 0 \) for
small $r$ because they do not have any significant effect on the central density of galaxy cluster. The baryonic component can be simply modelled by a constant density profile for small $r$ \cite{60}. It is significant only when $r \leq 0.1$ kpc.

Since the velocity profile is also time-dependent, Eq. (2) has to be solved numerically. Nevertheless, we can discuss some important features and behaviours of the resulting dark matter density profile by using some analytic approximations, especially for small $r$ regime.

In a typical galaxy, we have $M_{BH} \gg M_{DM}$ at small $r$. Therefore, the total mass profile is time and radius independent ($M(r,t) \approx M_{BH}$) near the galactic centre. In this regime, Eq. (2) gives

$$\rho(r,t) = \left[\frac{1}{\rho(r,0)} + \frac{<\sigma v_0 r^{\alpha/2}>}{(GM_{BH})^{\alpha/2}m} \right]^{-1} = \left( \frac{r}{\rho_s r_s} + K r^{\alpha/2} \right)^{-1}. \quad (3)$$

If $t$ is sufficiently large, $\rho(r,t)$ would follow $r^{-\alpha/2}$ at small $r$. When $M_{DM}$ dominates the mass profile, the total mass would be time and radius dependent. Nevertheless, if we neglect the time-evolution factor (it is a good approximation only near the core radius, and when $t$ is small), we can get an approximate analytic expression by using $M(r,t) \approx 2\pi \rho_s r_s r^2$ for $r < r_s$.

Eq. (2) gives

$$\rho(r,t) = \left[\frac{1}{\rho(r,0)} + \frac{<\sigma v_0 r^{\alpha/2}>}{(2\pi G \rho_s r_s r^{\alpha/2}m)} \right]^{-1} = \left( \frac{r}{\rho_s r_s} + K' r^{-\alpha/2} \right)^{-1}. \quad (4)$$

Moreover, since $M_{BH} = 0$ for small $r$ in most of the galaxy clusters, the baryonic component will dominate the mass profile for very small $r$ when most of the dark matter near the centre is annihilated ($t$ is sufficiently large). In this case, the total mass profile is just radius-dependent. We have

$$\rho(r,t) = \left[\frac{1}{\rho(r,0)} + \frac{<\sigma v_0 r^{\alpha/2}>}{(1.33 \pi G \rho_0 r^3)^{\alpha/2} m} \right]^{-1} = \left( \frac{r}{\rho_s r_s} + K r^{-\alpha} t \right)^{-1}, \quad (5)$$

where $\rho_0$ is the central baryon density. Since $\alpha > 0$, from Eqs. (4) and (5), we notice that $\rho(r,t)$ will rise again towards small $r$ for both galaxies and galaxy clusters.

From Eq. (4), we can see that there exists a turning point in the density when $K'r^{-\alpha/2}t \approx r/\rho_s r_s$ for dwarf and normal galaxies. Therefore, the core
radius of the structure $r_c$ should be the same order of magnitude as the position of the turning point. For $\alpha = 1$ and taking $\rho_s r_s = (58 - 144)M_\odot$ pc$^{-2}$ for dwarf and normal galaxy scale [55], we have $r_c \sim (10^{-4} - 10^{-3})(m/1$ GeV)$^{-2/3}$ kpc, which is too small to be the core radius. For $\alpha = 2$, we get $r_c \sim (0.1 - 1)(m/1$ GeV)$^{-1/2}$ kpc, which generally agrees with observations. This shows that our model is consistent only with $\alpha = 2$.

Let’s assume that $m = 1$ GeV, $f = 1$ and $\alpha = 2$. By solving Eq. (2) numerically with small time-steps, we can generate the time evolution of $\rho(r, t)$ for a typical dwarf galaxy, normal galaxy and galaxy cluster, respectively (see Figs. 1-3). All the parameters used are listed in Table 1. We use the empirical formula from [57] to model the corresponding mass of supermassive black hole in a typical dwarf galaxy and in a normal galaxy. We can notice from the figures that the density fluctuates at small $r$. This is consistent with the features obtained from the approximate analytic expressions in Eqs. (3)-(5). The inner slope is much shallower than the NFW profile. Secondly, neglecting the fluctuations of density for small $r$, we approximate the resultant density profile with a cored-NFW profile $\rho = \rho_c (1 + r/r_c)^{-1}(1 + r/r_s)^{-2}$ in each case. The core radii are $r_c \sim 1$ kpc and $r_c \sim 10$ kpc for galaxy and galaxy cluster respectively. The small core radii in dwarf galaxy and normal galaxy generally agree with observations [61, 34]. However, the structure of core in galaxy cluster is not obvious and the core radius is too small for us to observe. Current observational data from galaxy clusters for $r < 100$ kpc is highly uncertain. This may explain why core-like structures are not commonly found in galaxy clusters, even though they might exist.

The resulting density profile in our model is somewhat different from the result obtained from [31]. It is because the cross-section used in [31] is constant while the cross-section in this model is velocity-dependent. Fig. 4 shows a plot of the velocity profiles of dark matter particles in a typical galaxy and a typical dwarf galaxy at $t = 12$ Gyr. We can notice that the drop and rise in the resulting density profile strongly correlates with the velocity profile of dark matter. If there is no supermassive black hole or baryon in

| Structure        | $\rho_s$ (g cm$^{-3}$) | $r_s$ (kpc) | $M_{BH}$ ($M_\odot$) | $M$ ($M_\odot$) |
|------------------|------------------------|-------------|----------------------|-----------------|
| Dwarf galaxy     | $10^{-24}$             | 3.2         | $10^4$               | $10^{10}$       |
| Normal galaxy    | $3 \times 10^{-25}$    | 32          | $10^6$               | $10^{12}$       |
| Galaxy cluster   | $7 \times 10^{-26}$    | 500         | 0                    | $10^{15}$       |
the structure (i.e. $M_{BH} = 0$ and $M_B = 0$), the resulting density profile can be approximated by Eq. (4) only (density decreases towards the centre), and the shape of the density profile would be similar to the one shown in Fig. 3 because the effect of baryon in galaxy cluster is nearly negligible. If there is no Sommerfeld enhancement (i.e. $\alpha = 0$), from Eq. (4), the resulting central density profile would be constant ($\rho(t) \approx m/ <\sigma v_0> t$), which reduces to the result obtained from [31].

On the other hand, Fig. 5 shows that core-like structure disappears when $M_{BH}$ is large ($M_{BH} \sim 10^9 M_\odot$). Since large $M_{BH}$ usually correspond to large galactic mass [62], this suggests that core-like structures can be easily found in dwarf galaxies but not in large galaxies (with mass greater than $10^{13} M_\odot$).

If we take the local maximum turning point of the density profile being the core radius $r_c$, the mass profile at $r = r_c$ would be dominated by dark matter. At the turning point, from Eq. (4), we have $r_c \sim \sqrt{\rho_s r_s K t_0}$, where $t_0 \approx 12$ Gyr is the age of a galaxy. The central density is given by $\rho_c \approx \rho(r_c, t_0) \approx \rho_s r_s / r_c$. Therefore, we have $\rho_c r_c \approx \rho_s r_s$. As $\rho_s r_s \propto M^{0.2}$, our result is consistent with the empirical fits from observational data $\rho_c r_c \propto M^{0.2}$ [63]. Since this is a slow varying function of $M$, the product $\rho_c r_c \sim 100 M_\odot$ pc$^{-2}$ is nearly a constant for dwarf galaxies and normal galaxies, which agrees with the empirical fits using some cored-density profiles $\rho_c r_c = 141^{+82}_{-55} M_\odot$ pc$^{-2}$ [61, 64]. For $M \sim 10^{10} M_\odot$ in dwarf galaxies, our result gives $\rho_c \approx 0.07 M_\odot$ pc$^{-3}$, which matches the average central density ($\sim 0.01 - 0.1 M_\odot$ pc$^{-3}$) in dwarf galaxies [34].

Moreover, since many dark matter particles are annihilated, the central density become much smaller. This would provide a possible solution to solve the TBTF problem [11]. For $r < 250$ pc, the dark matter density is less than $10^{-23}$ g cm$^{-3}$ in a dwarf galaxy. Therefore, the rotational velocity would be less than 13 km/s, which can satisfactorily address the TBTF problem [65, 66].

4. Discussion

In this article, I show that the Sommerfeld-enhanced invisible dark matter annihilation is able to solve the core-cusp problem and obtain some general features that agree with observations. In order to produce the observed core size in galaxies, the value of $\alpha$ should be around 2, which means that the Sommerfeld boost is very close to resonance. Since the annihilation products
Figure 1: Time evolution of dark matter density profile in a dwarf galaxy. The thin solid line represents a cored-NFW density profile with $\rho_c = 5 \times 10^{-24}$ g cm$^{-3}$ and $r_c = 0.7$ kpc.

Figure 2: Time evolution of dark matter density profile in a normal galaxy. The thin solid line represents a cored-NFW density profile with $\rho_c = 1.5 \times 10^{-23}$ g cm$^{-3}$ and $r_c = 0.8$ kpc.
Figure 3: Time evolution of dark matter density profile in a galaxy cluster. The thin solid line represents a cored-NFW profile with $\rho_c = 5 \times 10^{-24}$ g cm$^{-3}$ and $r_c = 7$ kpc.

Figure 4: The circular velocity profile of dark matter in a typical galaxy (solid line) and a typical dwarf galaxy (dashed line) at $t = 12$ Gyr.
are invisible sterile neutrinos, the Sommerfeld-enhanced cross-section basically does not violate any observational bounds. This model can also give a satisfactory explanation why core-like structures commonly found in dwarf galaxies, but not in galaxy clusters. It is because the core size is too small to be observed in galaxy clusters. Moreover, this model just involves only two free parameters, the mass of supermassive black hole \( M_{BH} \), and the mass of a dark matter particle \( m \), which is assumed to be 1 GeV (the free parameter \( f \) can combine with \( m \) to become a single parameter \( f/m \)). In general, a smaller value of \( m \) would give a more obvious core-like structure and a lower central density. Since our result is consistent with the typical value of central density in galaxies, the value of \( m \) should be of the order 1 GeV.

Since a significant amount of dark matter would be annihilated, the dark matter content would be smaller in the inner region. Based on our calculations, more than 60% of the dark matter would be annihilated within 1 kpc. Surprisingly, this agrees with recent observations that the dark matter content is close to zero for the central part of galaxies [67]. Some observations also reveal that the dark matter content decreases with decreasing radius [68], which could also be explained by a higher rate of dark matter annihilation in the inner radius.

On the other hand, it is possible to have a very tiny baryonic branching ratio of the annihilation such that some photons or electron-positron pairs

---

Figure 5: The dark matter density profile \( \rho(r,t_0) \) in a normal galaxy with different supermassive black hole mass (solid line: \( M_{BH} = 10^6 M_\odot \); dashed line: \( M_{BH} = 10^8 M_\odot \)).
are directly created. However, the branching ratio should be less than $10^{-8}$ for these visible channels in order to satisfy the observational bounds \[33\].

In our model, although the inner slope of the density is much shallower than the NFW profile, the inner central density is not really a constant. However, due to observational uncertainties, the small fluctuations in the inner small region are usually smoothed out by using a constant density profile. If we have some good techniques in the future that can precisely probe the central density of dark matter, our model can be severely tested by three ways. First, the inner dark matter density is smaller than the NFW profile and goes like $1/r$ when $M_{BH}$ dominates the central mass. Second, there is a peak and trough in the dark matter density profile near the core radius. Third, our model predicts that core-like structures disappear in galaxies with large supermassive black holes. These features would be able to detect if we can precisely obtain the baryon density profile in galaxies or galaxy clusters in the future.

5. acknowledgements

This work is partially supported by a grant from the Hong Kong Institute of Education (Project No.:RG57/2015-2016R).

References

[1] J. F. Navarro, C. S. Frenk, S. D. M. White, Astrophys. J. 490 (1997) 493.

[2] R. A. C. Croft, D. H. Weinberg, M. Pettini, L. Hernquist, N. Katz, Astrophys. J. 520 (1999) 1.

[3] D. N. Spergel D. N. et al., Astrophys. J. Supp. 170 (2007) 377.

[4] E. Pointecouteau, M. Arnaud, G. W. Pratt, Astron. Astrophys. 435 (2005) 1.

[5] P. Salucci, Mon. Not. R. Astron. Soc. 320 (2001) L1.

[6] A. Borriello, P. Salucci, Mon. Not. R. Astron. Soc. 323 (2001) 285.

[7] E. Zackrisson, N. Bergvall, T. Marquart, G. Östlin G., Astron. Astrophys. 452 (2006) 857.
[8] W. J. G. de Blok W. J. G., Ad. Ast. 2010 (2010) 789293.

[9] M. Boylan-Kolchin, J. S. Bullock, M. Kaplinghat, Mon. Not. R. Astron. Soc. 415 (2011) L40.

[10] M. Boylan-Kolchin, J. S. Bullock, M. Kaplinghat, Mon. Not. R. Astron. Soc. 422 (2012) 1218.

[11] G. Ogiya, A. Burkert, Mon. Not. R. Astron. Soc. 446 (2015) 2363.

[12] P. Bode, J. P. Ostriker, N. Turok, Astrophys. J. 556 (2001) 93.

[13] Y.-J. Xue, X.-P. Wu, Astrophys. J. 549 (2001) L21.

[14] K. Abazajian, S. M. Koushiappas, Phys. Rev. D 74 (2006) 023527.

[15] M. Viel, J. Lesgourgues, M. G. Haehnelt, S. Matarrese, A. Riotto, Phys. Rev. Lett. 97 (2006) 071301.

[16] U. Seljak, A. Makarov, P. McDonald, H. Trac, Phys. Rev. Lett. 97 (2006) 191303.

[17] A. Boyarsky, J. Lesgourgues, O. Ruchayskiy, M. Viel, Phys. Rev. Lett. 102 (2009) 201304.

[18] A. V. Macciò, S. Paduroiu, D. Anderhalden, A. Schneider, B. Moore, Mon. Not. R. Astron. Soc. 424 (2012) 1105.

[19] D. N. Spergel, P. J. Steinhardt, Phys. Rev. Lett. 84 (2000) 3760.

[20] A. Burkert, Astrophys. J. 534 (2000) L143.

[21] N. Yoshida, V. Springel, S. D. M. White, Astrophys. J. 544 (2000) L87.

[22] S. W. Randall, M. Markevitch, D. Clowe, A. H. Gonzalez, M. Bradač, Astrophys. J. 679 (2008) 1173.

[23] J. Zavala, M. Vogelsberger, M. G. Walker, Mon. Not. R. Astron. Soc. 431 (2013) L20.

[24] A. Loeb, N. Weiner, Phys. Rev. Lett. 106 (2011) 171302.

[25] M. Vogelsberger, J. Zavala, A. Loeb, Mon. Not. R. Astron. Soc. 423 (2012) 3740.
[26] A. V. Macciò, G. Stinson, C. B. Brook, J. Wadsley, H. M. P. Couchman, S. Shen, B. K. Gibson, T. Quinn, Astrophys. J. 744 (2012) L9.

[27] J. Peñarrubia, A. Pontzen, M. G. Walker, S. E. Koposov, Astrophys. J. 759 (2012) L42.

[28] A. El-Zant, I. Shlosman, Y. Hoffman, Astrophys. J. 560 (2001) 636.

[29] C. Tonini, A. Lapi, P. Salucci, Astrophys. J. 649 (2006) 591.

[30] N. C. Amorisco, J. Zavala, T. J. L. de Boer, Astrophys. J. 782, L39 (2014).

[31] M. Kaplinghat, L. Knox, M. S. Turner, Phys. Rev. Lett. 85 (2000) 3335.

[32] J. F. Beacom, N. F. Bell, G. D. Mack, Phys. Rev. Lett. 99 (2007) 231301.

[33] M. S. Madhavacheril, N. Sehgal, T. R. Slatyer, Phys. Rev. D 89 (2014) 103508.

[34] J. Kormendy, K. C. Freeman, arXiv:1411.2170.

[35] A. Sommerfeld, Annalen der Physik 403 (1931) 257.

[36] L. F. Yang, J. Silk, A. S. Szalay, R. F. G. Wyse, B. Bozek, P. Madau, Phys. Rev. D 89 (2014) 063530.

[37] J.-J. Fan, P. Langacker, Journal of High Energy Physics 2012 (2012) 83.

[38] J. M. Cline, A. R. Frey, Phys. Lett. B 706 (2012) 384.

[39] R. Bernabei et al. [DAMA Collaboration], Eur. Phys. J. C 56 (2008) 333.

[40] C. E. Aalseth et al. [CoGeNT collaboration], Phys. Rev. Lett. 106 (2011) 131301.

[41] K. L. McDonald, Journal of High Energy Physics 2012 (2012) 145.

[42] G. Bertone, D. Hooper, J. Silk, Phys. Rept. 405 (2005) 279.

[43] J. B. Dent, S. Dutta, R. J. Scherrer, Phys. Lett. B 687 (2010) 275.

[44] J. Zavala, M. Vogelberger, S. White, Phys. Rev. D 81 (2010) 083502.
[45] S. Hannestad, T. Tram, JCAP 01 (2011) 016.
[46] C. Armendariz-Picon, J. T. Neelakanta, JCAP 03 (2014) 049.
[47] B. von Harling, K. Petraki, JCAP 12 (2014) 033.
[48] K. Griest, M. Kamionkowski, Phys. Rev. Lett. 64 (1990) 615.
[49] G. D. Mack, T. D. Jacques, J. F. Beacom, N. F. Bell, H. Yuksel, Phys. Rev. D 78 (2008) 063542.
[50] G. Rossi, C. Yeche, N. Palanque-Delabrouille, J. Lesgourgues, arXiv:1412.6763.
[51] J. M. Cornell, S. Profumo, W. Shepherd, Phys. Rev. D 88 (2013) 015027.
[52] J. L. Feng, M. Kaplinghat, H.-B. Yu, Phys. Rev. D 82 (2010) 083525.
[53] J. M. Cline, G. Dupuis, Z. Liu, W. Xue, Phys. Rev. D 91 (2015) 115010.
[54] J. Chen, Y.-F. Zhou, JCAP 04 (2013) 017.
[55] M. Schaller et al., arXiv:1409.8617.
[56] D. J. Croton, Mon. Not. R. Astron. Soc. 394 (2009) 1109.
[57] N. J. McConnell, C. P. Ma, Astrophys. J. 764 (2013) 184.
[58] A. E. Reines, J. E. Greene, M. Geha, Astrophys. J. 775 (2013) 116.
[59] A. C. Seth et al., Nature 513 (2014) 398.
[60] Y. Chen, T. H. Reiprich, H. Böhringer, Y. Ikebe, Y.-Y. Zhang, Astron. Astrophys. 466 (2007) 805.
[61] F. Donato, G. Gentile, P. Salucci, C. Frigerio Martins, Mon. Not. R. Astron. Soc. 397 (2009) 1169.
[62] K. Bandara, D. Crampton, L. Simard, Astrophys. J. 704 (2009) 1135.
[63] A. Del Popolo, V. F. Cardone, G. Belvedere, Mon. Not. R. Astron. Soc. 429 (2013) 1080.
[64] G. Gentile, B. Famaey, H. Zhao, P. Salucci, Nature 461 (2009) 627.
[65] M. G. Walker et al., Astrophys. J. 704 (2009) 1274.

[66] J. Wolf et al., Mon. Not. R. Astron. Soc. 406 (2010) 1220.

[67] F. Lelli, Galaxies 2 (2014) 292.

[68] S. A. Kassin, R. S. de Jong, B. J. Weiner, Astrophys. J. 643 (2006) 804.