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Puffy Dark Matter

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If dark matter has a finite size that is larger than its Compton wavelength, the corresponding self-interaction cross section decreases with the velocity. We investigate the implications of this Puffy Dark Matter for addressing the small-scale problems of the ΛCDM model. In particular, we show that the way the non-relativistic cross section varies with the velocity is largely independent of the dark matter internal structure when the range of the mediating force is very short. We present an explicit example in the context of a QCD-like theory of dark matter and show that low-threshold direct detection experiments have the potential to probe Puffy Dark Matter.

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Introduction. While the standard cosmology as described by the ΛCDM model has been well established thanks to the overwhelming amount of observational data, the particle nature of dark matter (DM) still eludes us. In this regard, the distribution of DM in astrophysical objects may provide a very important clue [1]. Actually, a handful of studies of the inner region of small-scale astrophysical halos claim that DM self-scatters with a cross section per unit of mass of $\sigma/m \gtrsim 1\,\mathrm{cm}^2/\mathrm{g}$ [2–4]. This is known as self-interacting dark matter (SIDM) and provides an appealing explanation to the seemingly mass deficit observed in objects such as dwarf galaxies when compared to the predictions of collisionless DM. See [5, 6] for recent reviews of these shortcomings of the ΛCDM model as well as for a discussion of alternative solutions such as those invoking other exotic DM candidates (e.g. [7–14]) or baryonic effects (e.g. [15–24]).

Yet, the aforementioned cross sections in small-scale objects are marginally consistent with observations of clusters of galaxies, in which DM moves relatively faster (see e.g. [25–27]). Thus, barring uncertainties, a velocity dependence of σ/m is preferred, with lower values at higher velocities. Due to this, point-like DM particles that self-scatter by means of a short-range interaction are often said to be disfavored because the corresponding cross section is nearly constant. In this context, mechanisms for obtaining a velocity-dependent σ/m include light mediators inducing a long-range interaction [28, 29], resonant SIDM [30], and processes involving inelastic scatterings [31–33].

In this work, we point out that, supposing that the DM particle has a finite size, $r_{\rm DM}$, the self-scattering cross section typically decreases for DM velocities larger than $(mr_{\rm DM})^{-1}$, even if the interaction associated with the scattering has a very short range. As is shown in Fig. 1, a momentum transfer much smaller than $r_{\rm DM}^{-1}$ is too small to measure the internal structure of the DM, so the latter acts as a point-like particle. On the other hand, when the momentum transfer becomes larger than $r_{\rm DM}^{-1}$, the internal structure of the particle is probed. As specified below, this can happen in such a way that the phase difference among the scattered waves leads to a suppression in the scattering cross section. This is indeed the desired velocity dependence of DM self-scattering. We will refer to this scenario as Puffy DM. We would like to remark that, beside the self-scattering effects, the fact that DM has a finite size leads to a very rich phenomenology, as has been explored for several concrete DM candidates (e.g. [34–51]).

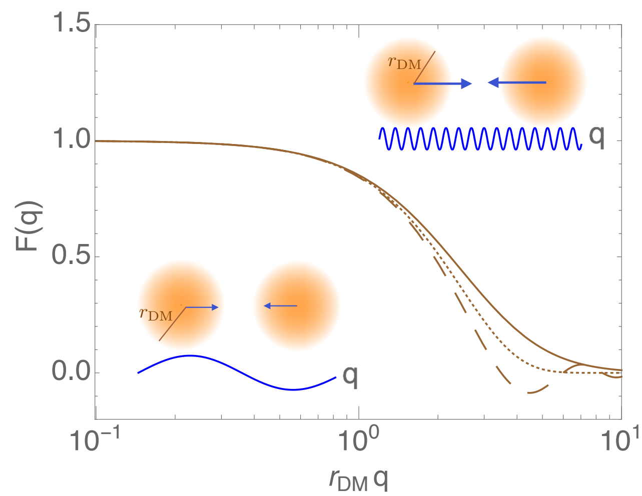

Figure 1: Form factors as a function of momentum-transfer q in units of the inverse root-mean-square radius $r_{\rm DM}$. Solid, dashed and dotted lines correspond to the dipole, the tophat and the Gaussian distributions (see Table I).

This manuscript is organized as follows. In Sec. II we elaborate further on the elastic scattering of finite-size DM particles. In Sec. III we discuss the implications in DM halos. In the following section we present a QCD-like model of Puffy DM. Finally we conclude in Sec. V.

II. Scattering of finite-size DM particles. Let us first consider the scattering of two finite-size objects, which –for simplicity– will be modeled as a collection of point-like constituents that coherently scatter by means

Shape	$\rho(r)$	$r_{\rm DM}$	$F(q)$
tophat	$\frac{3}{4\pi r_0^3}\theta(r_0-r)$	$2\sqrt{3}r_0$	$\frac{3(\sin(r_0q)-r_0q\cos(r_0q))}{r_0^3q^3}$
dipole	$\frac{e^{-r/r_0}}{8\pi r_0^3}$	$\sqrt{3/5}r_0$	$\frac{1}{(1+r_0^2q^2)^2}$
Gaussian	$\frac{1}{8r_0^3\pi^{3/2}}e^{-r^2/(4r_0^2)}$	$\sqrt{6}r_0$	$e^{-r_0^2q^2}$

Table I: Form factors for different density distributions.

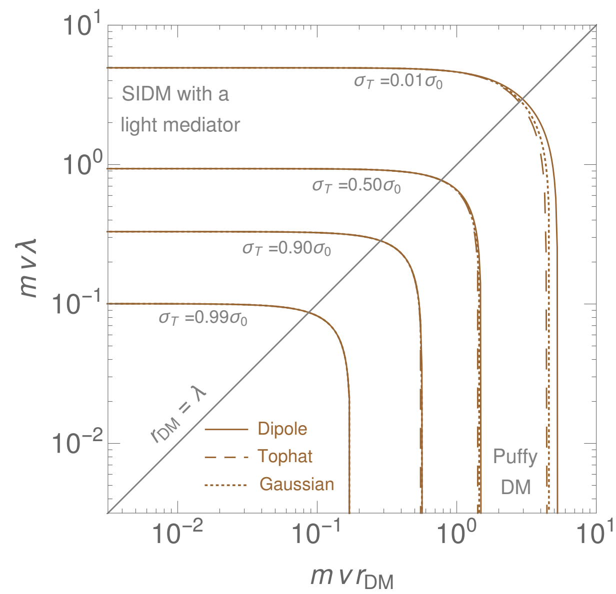

Figure 2: Transfer cross section as a function of the force range, λ, and the DM size, $r_{\rm DM}$, both in units of $1/mv$. Here σ_0 of Eq. (3) is assumed to be constant.

of a spin-independent Yukawa interaction. The corresponding charge density, $\rho(\vec{\mathbf{r}})$, characterizes the finite shape of the scattering object. We will also assume that the contribution of the binding force to the scattering rate is negligible. This is the case e.g. if such a force leads to a momentum-suppressed scattering amplitude. Then, the interaction Hamiltonian for two objects described by the density profiles $\rho_1(\vec{\mathbf{x}})$ and $\rho_2(\vec{\mathbf{y}})$ is

$$\begin{aligned} H_{int} &= \int d\vec{\mathbf{x}}d\vec{\mathbf{y}}\rho_1(\vec{\mathbf{x}})\frac{\alpha e^{-|\vec{\mathbf{x}}-\vec{\mathbf{y}}|/\lambda}}{|\vec{\mathbf{x}}-\vec{\mathbf{y}}|}\rho_2(\vec{\mathbf{y}}) \\ &= \int \frac{d\vec{\mathbf{q}}}{(2\pi)^3}F_1(\vec{\mathbf{q}})\frac{4\pi\alpha}{\vec{\mathbf{q}}^2+\lambda^{-2}}F_2(-\vec{\mathbf{q}})\,. \end{aligned} \qquad (1)$$

where λ is the range of the interaction, α is a coupling constant, and we have introduced the form factor $F_i(\vec{\mathbf{q}}) \equiv \int d\vec{\mathbf{r}}\, e^{i\vec{\mathbf{q}}\cdot\vec{\mathbf{r}}}\rho_i(\vec{\mathbf{r}})$. Hence, the center-of-mass differential cross section in the Born approximation is

$$\frac{d\sigma}{d\Omega} = S\left|F_1(\vec{\mathbf{q}})\frac{2\mu\alpha}{\vec{\mathbf{q}}^2+\lambda^{-2}}F_2(-\vec{\mathbf{q}}) \pm (\vec{\mathbf{q}}\to-\vec{\mathbf{q}})\right|^2, \qquad (2)$$

where μ is the reduced mass and $\vec{\mathbf{q}}$ is the momentum transfer. For identical (non-identical) particles, the second term must (not) be included and $S = 1/2\,(1)$.

An illustrative example is the electron scattering off larger objects. This is determined by a Coulomb interaction ($\lambda \to \infty$) with $F_e(\vec{\mathbf{q}}) = 1$. In this case, Eq. (2) gives the well-known Rutherford scattering formula, which can be used to infer the shape of finite-size objects. When applied to the proton, one finds a density distribution decreasing exponentially with a characteristic scale $r_0^{-2} = 0.71\,\mathrm{GeV}^2$ [52]. The latter is the dipole distribution (see Table I), generally expected from wavefunction solutions to various potential wells [53].

We apply now Eq. (2) to non-relativistic DM. Assuming that the DM particle is spherical, i.e. $F(\vec{\mathbf{q}}) = F(q)$, the S-wave differential cross section reads

$$\frac{d\sigma}{d\Omega} = \frac{\sigma_0}{8\pi}\left[\frac{F(q)^2}{1+\lambda^2q^2} + (\theta\to\pi-\theta)\right]^2_{q=mv\sin\theta/2}, \qquad (3)$$

where $\sigma_0 = 4\pi(m\alpha\lambda^2)^2$. Here θ and v are respectively the scattering angle and the relative velocity in the center-of-mass frame. While the exact form of $\rho(r)$ –and hence $F(q)$ in Eq. (3)– needs to be determined by solving for the wave function from the Schrödinger equation of the composite state, the differential cross section is not sensitive to the details of $\rho(r)$ as long as it is always positive (no screening) and it goes to zero sufficiently fast at large radii. In that case, the DM size –or more precisely– the root-mean-square radius

$$r_{\rm DM}^2 \equiv \int d\vec{\mathbf{r}}\,\rho(r)r^2 = -6\left.\frac{d^2F(q)}{dq^2}\right|_{q=0} \qquad (4)$$

is positive. Thus, $F(q)$ decreases for small momenta from $F(0) = \int d\vec{\mathbf{r}}\rho(r)$, which can be normalized to 1 without loss of generality. Fig. 1 illustrates this for the three representative distributions as listed in Table I. Together with Eq. 3, all this implies that the cross section is constant at low velocities and eventually approaches zero, even if the range of the interaction is extremely short.

III. DM scattering in astrophysical halos. Because of the form factor, for low velocities we expect isotropic scattering, whereas for larger velocities forward scattering is more probable. Due to this, the transfer cross section, $\sigma_T \equiv \int d\Omega(1-|\cos\theta|)d\sigma/d\Omega$, captures the self-interaction effects in DM halos better than σ (see e.g. [54]), and will be adopted below.

Fig. 2 illustrates the dependence of σ_T on the interaction range λ and the particle size $r_{\rm DM}$. As apparent from the plot, σ_T is largely independent of the exact expression for the form factor and therefore of $\rho(r)$. Furthermore, roughly speaking, the transfer cross section is constant for $mv \ll \min\{\lambda^{-1}, r_{\rm DM}^{-1}\}$, starts decreasing at $mv \sim \min\{\lambda^{-1}, r_{\rm DM}^{-1}\}$, and approximately scales as $1/v^4$

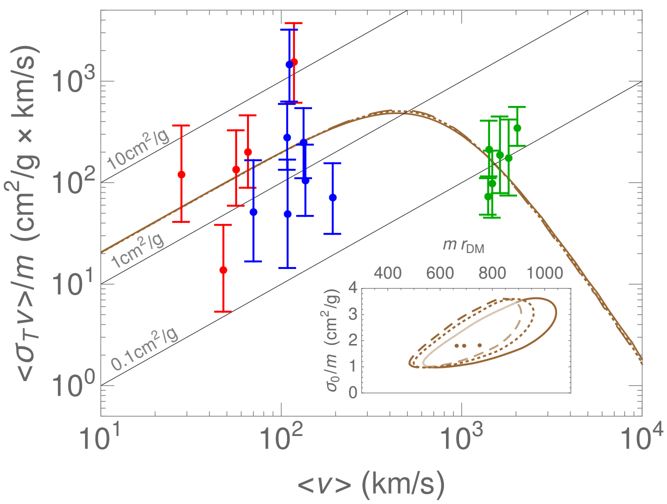

Figure 3: Velocity dependence of the transfer cross section of Puffy DM. Best-fit curves to data [27] for the dipole (solid), tophat (dashed) and the Gaussian (dotted) distributions in Table I. The inset shows the 95% C.L. contours together with the corresponding parameter sets of the main figure.

for $mv \gg \min\{\lambda^{-1}, r_{\rm DM}^{-1}\}$. See Appendix for details. Interestingly, when the range of the Yukawa force is much larger than the size of the DM, this precisely describes the Born regime of SIDM with a light mediator [29]. Furthermore, the figure shows that there is a one-to-one correspondence between the latter and the self-scattering of finite-size DM by a short-range force, both giving the same transfer cross section.

The DM relative velocity in astrophysical halos typically follows a Maxwell-Boltzmann distribution truncated at the corresponding escape velocity, $v_{\rm max}$. The velocity-averaged transfer cross section is then [68]

$$\langle \sigma_T v \rangle = \int_0^{v_{\rm max}} f(v)\sigma_T v dv\,, \quad f(v) = \frac{32 v^2 e^{-4v^2/\pi\langle v\rangle^2}}{\pi^2 \langle v\rangle^3}\,. \tag{5}$$

In the context of SIDM as a solution to the small-scale structure problems, a semi-analytical method has been proposed in [27] to infer, from observational data, the value of $\langle \sigma_T v\rangle/m$ for a given DM halo (see also [55]). This method was applied to five clusters from [56], seven low-surface-brightness spiral galaxies in [57] and six dwarf galaxies of the THINGS sample [58]. Fig. 3 shows these results respectively in green, blue and red. While these values should be taken with caution due to large uncertainties in the SIDM modeling of astrophysical objects (see e.g. [59]), the set of points is in agreement with observations from the Bullet Cluster giving $\sigma_T/m \lesssim 1.3\,{\rm cm}^2/{\rm g}$ at cluster scales [25, 26], which is one of the strongest constraints on DM self-interactions.

Postulating a DM finite size much larger than the range of the Yukawa force, i.e. $\lambda \ll r_{\rm DM}$, provides an excellent fit to the velocity-dependent cross section preferred by the galactic and cluster systems. The corresponding best-fit of Eq. (5) to the data above is shown in Fig. 3 for the dipole, the tophat and the Gaussian distributions, separately. As expected from the aforementioned remarks, there is almost no dependence on details of the form factors even though they correspond to substantially different density distributions. The figure also shows that, in order to have the right velocity dependence, the DM size needs to be hundreds of times larger than the Compton wavelength. This explains the name Puffy DM.

If the Yukawa force is associated to a mediator ρ, requiring $\lambda = 1/m_\rho \ll r_{\rm DM}$ implies $m_\rho \gg 10^{-3} m$. This shows that the mediator can be lighter than the DM and still the velocity dependence is determined by the DM size. Likewise, if we impose $\alpha \lesssim m_\rho/m$ as required in the Born expansion, $\sigma_0/m = 4\pi(m\alpha\lambda^2)^2/m \sim 1\,{\rm cm}^2/{\rm g}$ leads to $m \lesssim 20\,{\rm GeV}$.

IV. A model of Puffy DM. Here we only sketch a possible realization of Puffy DM while details will be discussed elsewhere. It is a QCD-like confining theory with N_c colors and two flavors of quarks: one "charm quark" much heavier than the confining scale Λ and one nearly massless "down quark". They respectively have charges $+2/3$ and $-1/3$ under a dark $U(1)_D$ gauge group. This is associated with a massive "dark photon" γ_D, which can act as a portal to the Standard Model (SM) by means of the kinetic mixing between the $U(1)_D$ group and the SM hypercharge. There are no dark weak interactions. We assume there is an asymmetry so that anti-charm quarks are annihilated while the remaining charm quarks end up in the baryonic $\Sigma_c(cdd)$ state. The latter interacts by exchanging the pseudo-scalar $\eta(d\bar{d})$ and the vector $\rho(d\bar{d})$, which lead to attractive and repulsive forces respectively.

On the one hand, it is likely that the η-exchange dominates binding Σ_c baryons into nuclei because its range is larger given that the η mass is due to the anomaly and hence suppressed as $m_\eta \sim \Lambda/\sqrt{N_c}$, as opposed to the ρ mesons for which $m_\rho \sim \Lambda$. In view of this, in the following we assume the typical mass number is $10 \lesssim A \lesssim 100$. On the other hand, the nucleus-nucleus scattering is dominated by the exchange of ρ mesons because the latter are essentially massive gauge bosons coupled to d-number $(A/2)$ giving rise to coherent spin-independent scattering, while the η-exchange induces a spin-dependent momentum-suppressed scattering. Therefore, the range of the scattering force Λ^{-1} is shorter than the size of the nuclei $r_{\rm DM} \sim A^{1/3} m_\eta^{-1} \sim A^{1/3}\Lambda^{-1}\sqrt{N_c}$. As a result, this model is a realization of Puffy DM.

For instance, parameters such as $N_c = 3$, $A \sim 10$, $m_c \sim m_{\Sigma_c} \sim 1\,{\rm GeV}$, $r_{\rm DM}^{-1} \sim 15\,{\rm MeV}$, $m_\eta \sim 20\,{\rm MeV}$, $\Lambda \sim m_\rho \sim 30\,{\rm MeV}$ and $\alpha \sim m_\rho/m$ realize the desired self-scattering cross section and its velocity dependence.

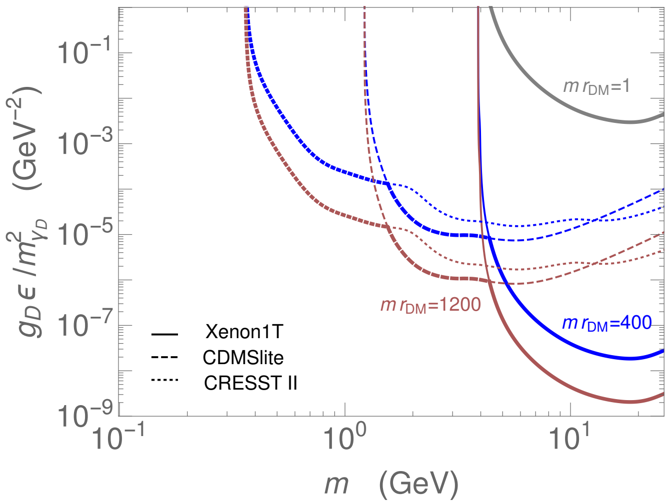

Figure 4: Direct detection bounds on our QCD-like theory of Puffy DM from Xenon1T [64], CMDSlite[65] and CRESST II [66]. See text for details.

We take $m_{\gamma_D} < m_\eta/2$ so that η decays into $\gamma_D\gamma_D$ from the anomaly [69]. Then the size of the kinetic mixing is either (A) $10^{-5} \lesssim \epsilon \lesssim 10^{-3}$ or (B) $\epsilon \ll 10^{-10}$ to satisfy beam-dump experimental data and supernova observations [60, 61]. In the cosmological history, presumably much of the entropy in this sector ends up in a thermally populated gas of η mesons. These decay via $\eta \to \gamma_D\gamma_D \to 2(e^+e^-)$ before Big-Bang Nucleosynthesis (BBN) for the range (A) if $g_D \gtrsim 10^{-9}$. On the other hand, the direct detection forces $g_D \lesssim 10^{-6}$ for a 10 GeV DM particle.

Even though the nucleus is $U(1)_D$ neutral, it has a finite charge radius similar to the neutron, which approximately equals the DM size $r_{\rm DM}$. Therefore the kinetic mixing between our photon and the dark photon leads to nuclear recoils. We estimate the corresponding current direct-detection limits by implementing such a recoil spectrum in DDCalc [62, 63]. The results are shown in Fig. 4 for various choices of the charge radius and $1\,{\rm MeV} \ll m_{\gamma_D} \ll m_Z$.

For the parameter range (B), the dark sector decouples from the Standard Model early, and hence it may be much cooler and the additional entropy ejection is limited. This case needs to be studied separately.

V. Conclusions. We have shown that if DM is a composite state with a size hundreds of times larger than its Compton wavelength, the corresponding self-interaction cross section varies with velocity in a way that is largely independent of its internal structure. For cross sections larger than $1\,{\rm cm}^2/{\rm g}$ at $v \to 0$, this provides a solution to the problems of the ΛCDM model in small-scale astrophysical objects while still being in agreement with cluster observations. An important aspect of this scenario is that it does not require a long-range force mediating DM self-scatterings. A QCD-like theory where DM is a dark nucleon has been used to illustrate our results, which are nevertheless general and can be applied to a broader range of theories. For this reason, we believe Puffy DM opens up a new avenue for SIDM model-building.

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Appendix: The transfer cross section

The transfer cross section for DM scattering is

$$\sigma_T = \frac{\sigma_0}{8\pi}\int d\Omega\,(1-|\cos\theta|)\left[\frac{F(q)^2}{1+\lambda^2 q^2} + (\theta\to\pi-\theta)\right]^2 = \int_0^{\frac{(mv)^2}{2}} \left[\frac{F(q)^2}{1+\lambda^2 q^2} + \left(q^2 \to (mv)^2 - q^2\right)\right]^2 \frac{2\sigma_0 q^2 dq^2}{(mv)^4}. \quad (6)$$

Here we focus on the Puffy DM, where $r_{\rm DM}^{-1} \ll \lambda^{-1}$. On the one hand, taking the low velocity limit, $mv \ll r_{\rm DM}^{-1}$, the factor in the square bracket approaches 2 and thus $\sigma_T \to \sigma_0$ at $v \to 0$. On the other hand, for $mv \gg r_{\rm DM}^{-1}$, $F(q)$ is suppressed at $q \gg r_{\rm DM}^{-1}$ so that the result of the integral is insensitive to its upper limit. The integral is not sensitive to λ either, because for any $q \gtrsim \lambda^{-1}$ there is always $q \gg r_{\rm DM}^{-1}$. Taking this into account allows us to approximate σ_T by

$$\sigma_T \simeq \frac{2\sigma_0}{m^4 v^4}\int_0^\infty dq^2 q^2 F(q)^4 \simeq \frac{\sigma_0}{(c\,mvr_{\rm DM})^4}\,, \quad (7)$$

with $c = 0.23$, 3.9, 0.97 for the tophat, the dipole, and the Gaussian distributions, respectively. Therefore, σ_T scales as $1/v^4$ at $mv \gg r_{\rm DM}^{-1}$. The behavior derived here is different from that of effective range theories [67], since the latter applies to each partial wave of the scattering cross section, while our result applies to the total transfer cross section. Note that at very large mv incoherent scattering starts playing a role. Nevertheless, its contribution is much smaller than σ_T, and is therefore neglected for simplicity.

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[68] Here we will assume that $v_{\max} \to \infty$ since the integral converges quite fast due to the Boltzmann factor.
[69] For slightly heavier dark photons the process

$\eta\eta \to \gamma_D\gamma_D$ can sufficiently reduce the abundance of η mesons. If a dark axion, a, exists, the annihilation $\eta\eta \to \eta a$ can play the same role.