We show that the mere observation of the first stars (Pop III stars) in the universe can be used to place tight constraints on the strength of the interaction between dark matter and regular, baryonic matter. We apply this technique to a candidate Pop III stellar complex discovered with the Hubble Space Telescope at $z \sim 7$ and find bounds that are competitive with, or even stronger than, current direct detection experiments, such as XENON1T, for dark matter particles with mass ($m_X$) larger than about 100 GeV. We also show that the discovery of sufficiently massive Pop III stars could be used to bypass the main limitations of direct detection experiments: the neutrino background to which they will be soon sensitive.

Keywords: dark matter; dark matter capture; stars

One of the most intriguing open questions in Physics today is the nature of Dark Matter (DM). Its existence has been inferred via the gravitational effects it has from the smallest scales, in the Cosmic Microwave Background (CMB) radiation [1–4], to intermediate, galactic [5] and cluster scales [6]. The complex large scale structures and sub-structures dark matter forms during its gravitational collapse around the potential wells provided by the primordial density fluctuations can be mapped using gravitational lensing on galactic cluster [7] and cosmological [8–10] scales. Over the past few decades a standard, concordance cosmological model has emerged as a leading candidate that best explains all available cosmological data: the Λ-CDM model. About 27% of the energy budget of the Universe today is in the form of DM, whereas regular, baryonic, matter only amounts to roughly 5%. The other 68% is thought to be comprised of Dark Energy: a uniform, negative pressure fluid responsible for the current accelerated expansion of the Universe. One should note that recently there are hints of tensions between data and the concordance model [11][12]. Most notably the discrepancy in the current expansion rate of the universe, $H_0$, as inferred from late and early Universe probes [13][14]. On small scales, there are a number of challenges that the Λ-CDM paradigm faces [15].

Dark Matter detection. Currently, there are three broad strategies in the hunt for dark matter: production of DM particles in accelerators, direct, and indirect detection. Each of these exploits the various possible interaction channels between dark matter and baryonic matter. The Large Hadron Collider (LHC) has not yet found any evidence of particles outside of the standard model; as such, the minimum mass of any dark super-symmetric DM particle candidate has been pushed to higher and higher values. Indirect detection experiments rely on the possible self-annihilations or decay of dark matter particles, whenever DM densities are high. The nearest such site is the center of our own galaxy. An antiproton and a gamma-ray excess compared to known backgrounds have been found in Alpha Magnetic Spectrometer (AMS) and Fermi satellite data, respectively. Intriguingly, both of those excesses could be fit with a $\sim 60$ GeV DM particle self annihilating [16–18]. Alternatively, there are astrophysical explanations for those excesses [19][20].

Direct detection experiments exploit the small amount of energy a dark matter particle deposits as it collides with atomic nuclei [22][23]. As such, they are extremely challenging; moreover, shielding from overwhelming cosmic ray backgrounds requires performing the experiments deep underground. So far, DAMA/LIBRA is the only group that reports a signal consistent with DM detection [24][25]. Unfortunately, this signal, reported since 1998, has not been confirmed by any other laboratory. In lack of a clear detection signal, direct detection experiments are constraining the allowed strength of the interaction between dark matter and baryonic matter. As they become more and more sensitive, their detectors will be swamped with signals from neutrinos, which cannot be disambiguated from any possible dark matter signal. At that stage, if no clear DM signal identification has been made, new detection strategies will need
to be implemented. For reviews on dark matter and its detection status see Refs. 27 [52].

In this letter we propose a novel method of constraining the dark matter proton scattering cross section using Pop III stars, applicable when DM self annihilates. Using this formalism for the candidate Pop III system at redshift \( z \sim 7 \), found in the Hubble Space Telescope (HST) data [33], we obtain some of the most stringent bounds to-date, competitive with XENON1T. Moreover, we show that the upcoming James Webb Space Telescope (JWST), and its potential for discovering massive Pop III stars, could be used to probe below the neutrino scope (JWST), and its potential for discovering massive Pop III stars, applicable when DM self annihilates. For reviews on dark matter and its metric signatures compared to Pop III stars [45, 46]. In this work we assume that Dark Stars and Pop III stars are not mutually exclusive, and that at least some of the first stars are Pop III stars.

Any astrophysical object can accrete dark matter at its core via a phenomenon called capture [47–49]. Pop III stars formed via the gravitational collapse of zero metallicity, primordial baryonic gas clouds that contain pristine H and He from big bang nucleosynthesis. This happened at high redshifts (\( z \sim 10–50 \)) [7] at the center of DM mini-halos (\( M_{\text{h halo}} \sim 10^6 M_\odot \)), in very DM-rich environments. Using hydrodynamical simulations, the following picture emerges: typically one or just a few Pop III stars form per mini-halo, with masses up to \( \sim 1000 M_\odot \), powered by \( H \) fusion. For reviews on the formation of the first stars, see Refs. 35 [42]. Under certain conditions [43], DM heating during the formation of the first stars leads to objects powered by DM annihilations, Dark Stars (DS). Those hypothetical objects can grow to be supermassive [41], and have different photometric signatures compared to Pop III stars [35] [40]. In this work we assume that Dark Stars and Pop III stars are not mutually exclusive, and that at least some of the first stars are Pop III stars.

Any astrophysical object can accrete dark matter at its core via a phenomenon called capture [47–49]. Pop III stars, forming in a DM-rich environment, are particularly good probes of this phenomenon. Refs. [50] [51] study this for weakly interacting (WIMP) dark matter that gets captured by at most one collision (single-scattering) with nuclei inside Pop III stars. Using the recently developed multiscattering capture formalism [52] two of the authors of this letter investigated the capture of superheavy (\( m_X \gtrsim 10^8 \) GeV) dark matter by Pop III stars [54], and found that heating from dark matter annihilations leads to an upper bound on the Pop III masses. In this letter, we show how the mere observation of a Pop III star of any given mass can be used to constrain the DM-proton scattering cross section. Below we briefly summarize our method.

Any star of a given mass can never shine brighter than the Eddington luminosity (\( L_{\text{Edd}} \)):

\[
L_{\text{nuc}}(M_\star) + L_{\text{cap}}(M_\star; \text{DM params}) \leq L_{\text{Edd}}(M_\star) \quad (1)
\]

\( L_{\text{nuc}} \) represents the heating generated by the hydrogen fusion at the core of the star, whereas \( L_{\text{cap}} \) is the heating due to captured dark matter annihilations, which depends both on stellar [7] and DM parameters. Most importantly, it is sensitive to the DM-proton scattering cross section (\( \sigma \)). This ultimately allows us to place bounds on \( \sigma \), if all other parameters are measured or constrained; conversely, we can use the current bounds on \( \sigma \) from XENON1T to predict what the maximum mass of a Pop III star would be, as an effect of captured DM heating. Below we briefly summarize the technical details necessary for constraining \( \sigma \), or predicting an upper mass on \( M_\star \). For more details, please consult the companion paper [55].

DM particles crossing a star with radius \( R_\star \) can, via collisions with nuclei, lose enough energy to become trapped by the gravitational field of the star. This happens when the DM particle velocity falls below the escape velocity (\( v_{\text{esc}} \)) of the star. The capture rate is given by [52]:

\[
C_{\text{tot}} = \sum_{N=1}^{\infty} C_N = \sum_{N=1}^{\infty} \pi R_\star^2 \times \frac{\rho_X}{m_X} \int_0^{u_{\text{max}}} \frac{f(u)du}{(u^2 + v_{\text{esc}}^2)} \times \frac{p_N(\tau)}{p_N(\tau)} \quad (2)
\]

The probability a DM particle will collide exactly \( N \) times as it crosses the star has the following closed form [53]: \( p_N(\tau) = \frac{1}{\tau} \left( N + 1 - \frac{\Gamma(N+2,\tau)}{\Gamma(N+2)} \right) \), where \( \Gamma(a,b) \) is the incomplete gamma function. The optical depth is defined as: \( \tau = 2R_\star \sigma n_T \), where \( n_T \) is the average number density of nuclei inside the star. Throughout, \( \rho_X \) represents the DM density. DM particles with velocity \( u \) (measured infinitely far from the star) greater than \( u_{\text{max}} = v_{\text{esc}} \left( (1 - \beta_+ / 2)^{-N} - 1 \right)^{1/2} \) will not be captured after \( N \) collisions, since they are too fast to be slowed below \( v_{\text{esc}} \). Here, \( \beta_+ \equiv 4m_Xn_X/(m + m_X)^2 \), with \( m \) being the mass of the target nucleus. For this reason, we only integrate the velocity distribution up to the \( u_{\text{max}} \) cutoff. This amounts to only a part of the DM particle flux crossing the star being captured. The key point is

\footnotesize
1 Sometimes at \( z \) as low as 7 [7] 2 See also [53]. 3 Homology relations relate \( R_\star \) with \( M_\star \) [55]

that, as expected, the capture rate depends on the scattering cross section (via $r$). In \[53\] we presented closed form analytical expressions for $C_{\text{tot}}$, obtained assuming a Maxwell-Boltzmann distribution $f(u)$. For details of the calculation see the companion paper \[55\].

After being captured by a Pop III star, DM particles enter an equilibrium regime, where the capture rate ($C_{\text{tot}}$) equals twice the annihilation rate, thus the total number of DM particles remains constant \[54\]. Remarkably, in this regime the capture rate controls the heating due to DM annihilations: $L_{\text{cap}} = 2fC_{\text{tot}}m_X$. Henceforth, $f$ represents the fraction of the annihilation energy that is deposited inside the star, for which we assume, following \[32\], a value of 2/3. For details on how $L_{\text{nucl}}$ or $C_{\text{tot}}$ are calculated, see the companion paper \[55\]. In the next few paragraphs we summarize those results.

In practice, we calculate $C_{\text{tot}}$ numerically by summing the $C_N$’s of Eq. 2 up to a cutoff, $N_{\text{cut}}$, when the sum has converged. We find that $N_{\text{cut}} \propto r$. The capture rate in the multiscattering capture regime, i.e. $r \gg 1$, has two different scalings. First, note that if $u_{\text{max}} \to \infty$, then the capture rate becomes: (star cross sectional area)×(Total Flux), i.e. the number of DM particles crossing the star. This is obviously insensitive to $r$. Therefore, the constraining power of our method is lost in the region of parameter space corresponding to this scenario. The $u_{\text{max}} \to \infty$ condition is equivalent to the $k\tau \gg 1$ (i.e. Region II of Fig. 1), where $k = \frac{3v^2}{2} \leq m_X^2 M_* \left[\frac{2}{m_X \bar{v}}\right] \left[1 - \frac{1 - e^{-B^2}}{B^2}\right]$. (3)

$\bar{v}$ is the dispersion velocity of the DM distribution, $M_*$ and $R_*$ being the mass and the radius of the star, respectively. Current XENON1T bounds on $\sigma$ guarantee that when $m_X \lesssim 10^{10}$GeV, Pop III stars will capture DM in the single scattering regime, since $r \ll 1$ for that mass range. For this case, the capture rate has been calculated first by \[49\]. Up to numerical constants, it scales as:

$$C_1 \approx \rho_X \sigma \frac{M_*^3}{m_X^3 \bar{v}^3 R_*^2} \left[1 - \frac{1 - e^{-B^2}}{B^2}\right].$$ (4)

Here $B = \frac{3 \nu_{\text{esc}}^2}{2} \left[\frac{m_X}{m_\nu}\right] \frac{4\pi}{(\mu+1)^2}$, with $\mu \equiv \frac{m_X}{m_\nu}$, i.e. the ratio between the DM particle mass and the target nucleus mass. When $B \ll 1$, i.e. at the higher $M_\nu$ end (Region IV), remarkably, we recover the same scaling in the multiscattering regime given by Eq. 3. We reached the same conclusion by taking the $r \ll 1$ limit (i.e. single scattering) of the total capture rate for the multiscatter formalism \[55\]. In fact, in our numerical work we use the multiscatter formalism exclusively, since it naturally incorporates the single scattering limit. For $B \gg 1$, i.e. at lower $m_X$ (Region III), since the term in the square bracket can be approximated with one, we have $C_1 \approx \rho_X \sigma \frac{M_*^3}{m_X^3 \bar{v}^3 R_*^2}$, as found by \[54\], and confirmed numerically in the companion paper \[55\], by using the multiscatter formalism.

In view of homology relations \[55\], the capture rates, and implicitly $L_{\text{cap}}$, depend only on the following set of parameters: $\rho_X, m_X, \sigma, \bar{v}, M_*$. For the nuclear luminosity, we find the following interpolating function:

$$L_{\text{nucl}} \approx 10^{\frac{2\log(3.71 \times 10^3 L_{\text{XENON1T}})}{1 + \exp(-0.85 x - 1.955)}} \times 2 \times 10^{606 + x} \text{ erg/s}$$ (5)

where $x = \frac{M_\nu}{M_*}$ and $L_\odot \equiv 3.846 \times 10^{33}$ erg/s. As expected, this logistic fit function transitions between $L_{\text{nucl}} \approx M_*^3$ for intermediate mass stars, to $L_{\text{nucl}} \approx M_\nu$, for $M_* \gtrsim 1000M_\odot$. For the Eddington luminosity, assuming BBN composition of Pop III stars, we find: $L_{\text{Edd}} \approx 3.71 \times 10^4 (M_\nu/M_\odot) L_\odot$.

We end this section with a figure summarizing our method of constraining the scattering cross section, assuming a Pop III star of a given mass is observed. In

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**FIG. 1.** Exclusion limits in $\sigma$ vs. $m_X$ parameter space placed using a hypothetical 100$M_\odot$ Pop III star for two different assumed DM densities (blue vs red lines). Note the various regions separated by the line of $r = 1$ (multiscatter vs single scatter capture), $k = 1$ and $k\tau = 1$, with $k \equiv \frac{3\nu_{\text{esc}}^2}{2} \left[\frac{m_X}{m_\nu}\right] \frac{4\pi}{(\mu+1)^2}$. Region II corresponds to the scenario where the cross section would be so high, that essentially all DM particles crossing the star would be captured, rendering the capture rate insensitive to $\sigma$, and therefore limiting our method. Note that the bounds placed scale inversely with the ambient DM density. For a 100$M_\odot$ Pop III star we find that the exclusion bounds at $m_X \gtrsim 10^6$GeV would precisely match the current XENON1T bounds if $\rho_X = \rho_{\text{XENON1T}} \sim 10^{15}$ GeVcm$^{-3}$. For more massive stars this value would be lower. Additionally, for a given Pop III mass, the ambient DM density greater than its corresponding $\rho_{\text{XENON1T}}$ would lead to bounds that are deeper than current XENON1T ones.

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Fig. 1 we plot the exclusion bounds obtained using a
hypothetical $100M_\odot$ Pop III star for two critically important ambient DM densities. First, the density labeled
\[ \rho_{\text{crit}} = \frac{L_{\text{cap}}(M_\odot) - L_{\text{cap}}(M_\star)}{\frac{1}{2} \sqrt{\frac{2 \pi^2}{3}} \sigma^2 R_\star^2} \].

For densities lower than $\rho_{\text{crit}}$, our method can no longer be used to constrain $\sigma$ vs $m_X$, as this corresponds to the regime where all DM particles crossing the star are captured, independent of $\sigma$ (Region II of Fig. 1). We point out that for Pop III stars, this region has already been excluded by XENON1T, so in practice it will not be a limitation of our method. The other DM density considered in Fig. 1 $\rho_X$, will lead to bounds that precisely overlap, at large $m_X$, with the XENON1T current exclusion limits. Our exclusion limits, at large $m_X$, will be bound by $\sigma \sim m_X$ lines. This can be easily understood from the fact that $L_{\text{cap}} \sim \frac{\sigma}{m_X}$ (see Eq. (3)) in the multiscattering regime (Region I in Fig. 1). At intermediate $m_X$, we enter the single scattering regime. As pointed out in the discussion of Eq. (4), there are two different scalings of $C_{\text{tot}}$ in the single scattering regime, corresponding to the different behaviours of our exclusion limits: $\sigma \sim m_X^4$ (in Region III) vs. $\sigma \sim m_X^2$ (in Region IV).

We sum up our method: by using the sub-Eddington condition (Eq. (1)) we can find an upper bound on $M_\star$ for Pop III stars, when $\sigma$ is constrained via direct detection experiments. Conversely, once a Pop III star with a given mass is identified, we can use that information to place constraints on $\sigma$ as a function of $m_X$. In the next section we discuss those results in detail.

Results. In Fig. 2 we present upper bounds on Pop III stellar masses obtained by imposing the sub-Eddington condition, Eq. (1). For $\dot{v}$ we have assumed a fiducial value of 10 km/s, appropriate for the minihalos hosting Pop III stars. As expected, for a given $m_X$, an increase in the ambient DM density, $\rho_X$, leads to tighter bounds. We also note that at a fixed $\rho_X$, the bounds are only sensitive to $m_X$ at the lowest part of the mass end. For high $m_X$, capture happens in the multiscattering regime, and therefore $C_{\text{tot}} \propto \sigma/m_X^4$ (see Eq. (3)). For the DM particle mass range considered in this paper, current upper bounds on $\sigma$ from direct detection experiments scale linearly with $m_X$. Therefore, the upper bound on $L_{\text{cap}} \sim m_X C_{\text{tot}} \propto m_X^5$, i.e. is insensitive to $m_X$. At lower DM mass, in the regime when $C_{\text{tot}} \sim \sigma/m_X$, we have $L_{\text{cap}} \propto m_X$. Therefore the upper bound on $M_\star$ will increase as we decrease $m_X$, a trend that can be seen in Fig. 2. Note that captured DM annihilations can lead to maximum Pop III stellar masses as low as one stellar mass, for sufficiently high ambient DM densities!

The most exciting application of our method is the possibility to constrain the DM-proton scattering cross section, once we know the mass of any Pop III star (see Fig. 1). The main limitation comes from the uncertainty in the determination of $\rho_X$, the ambient DM density at the location of the star, since there is no possible dynamical determination one can make for halos that distant.

We assume an adiabatically contracted Navarro-Frenk-White (NFW) DM profile for the host minihalo. As the baryonic protostellar cloud cools and collapses, it will modify the initial DM density profile by enhancing densities in the inner regions of the halo. This is simply a response of the DM orbits to an increase in the gravitational potential. The Adiabatic Contraction (AC) formalism can be used to estimate this DM density enhancement, using the simplifying assumption of the existence of adiabatic invariants for DM particles inside a halo. Results from numerical simulations are in good agreement with those obtained via the adiabatic contraction formalism, especially for high redshift halos, such as those where Pop III stars form, since baryonic feedback effects are not important in this case. In older galaxies, Active Galactic Nuclei, or radiative feedback from very massive stars, can lead to a suppression of the infall of baryons, and therefore a suppression of the enhancement of the DM densities.

We point out that even for direct detection experiments, for which rotation curves of the Milky-Way galaxy can be used to determine radial distribution of the total mass, the DM density and velocity distribution in the solar neighbourhood are still uncertain to the level of 10%. Typically, one assumes a local DM density of 0.3 GeV cm$^{-3}$, as per the Standard Halo Model (SHM). However, recent simulations show that this value could in fact be larger, by $\sim 10\%$, i.e. $\rho_0 = 0.33$ GeV cm$^{-3}$. It is remarkable that this value is also favoured by the Gaia DR2 data, as shown in Ref. [65]. Moreover, Ref. [65] demonstrates that Milky-Way rotation curve data tends to prefer the physically motivated contracted NFW halo, which can be seen as direct experimental evidence of the compression of dark matter.
densities due to baryonic infall.

In [55], we use the adiabatic contraction formalism to calculate the DM ambient densities relevant to DM capture by Pop III stars. We show that the DM densities at the edge of the collapsing baryonic core during the formation of Pop III stars can attain values as high as $10^{19}$ GeV cm$^{-3}$, assuming the adiabatic compression operates until the protostellar core reaches hydrostatic equilibrium, at a hydrogen number density of $n \sim 10^{22}$ cm$^{-3}$. Up to factors of order unity, this estimate holds for a large variety of concentration parameters for the initial NFW profile ($c \sim 1$ to $c \sim 10$) and for redshifts ranging from $z \sim 20$ to $z \sim 5$. If, however, we adopt a more conservative approach, and assume that adiabatic compression stops operating earlier than the formation of the proto Pop III star, we get a lower value for the ambient DM density. For instance, for $n \sim 10^{16}$ cm$^{-3}$, we estimate the DM density at the edge of the baryonic core to be $\sim 10^{15}$ GeV cm$^{-3}$. At any rate, for the conservative approach, the value quoted is an underestimate of the actual density at the boundary of Pop III stars, which is the relevant parameter. This is because a typical Pop III star has a radius $R_* \sim R_{\odot}$, whereas the edge of the baryonic core corresponds to $r_B \gg R_*$, where the DM density is lower. We note these estimates are robust against changes in the initial DM density profile [66], and numerical simulations give very similar results [56]. Simulations are resolution limited, and currently they can probe the DM density only from $\sim 10^{-2}$ pc outward. For DM densities closer to the center of the DM halo, for now, we have to rely on the adiabatic compression approximation.

In Fig. 3 we present our bounds on the DM-proton scattering cross section, contrasted against the current, deepest available exclusion limits from direct detection experiments. The system we used as a DM probe was found in the MUSE deep-lensed field with the Hubble Space Telescope (HST) by [33]. They show that the Ly$\alpha$ emission from this $z \sim 7$ system can be modeled by Pop III stars with masses ranging from 100$M_{\odot}$ to 1000$M_{\odot}$. If confirmed with JWST, this would be the first discovery of a zero metallicity Pop III stellar system! Note that our bounds are the same for spin-dependent (SD) or spin-independent (SI) interactions. This is in contrast to direct detection experiments on Earth, for which the SD bounds are typically weaker by about five orders of magnitude. Even for the conservative $\rho_X \sim 10^{14}$ GeV cm$^{-3}$, all of our exclusion limits rule out a large swath of parameter space for SD proton-DN cross sections that has yet to be explored by direct detection experiments.

For the case of a 1000$M_{\odot}$ Pop III star, even our more conservative bounds are competitive with the XENON1T SI limits, as one can see from the upper green dashed line of Fig. 3. Finally, we point out that for a Pop III star of any given mass, there is a corresponding DM density ($\rho_{X,\text{mf}}$) for which the mere existence of the star in question will rule out DM-proton cross sections all the way down to the neutrino floor. At the same stellar mass (or $\rho_X$), a higher $\rho_X$ (or $M_*$) implies probing below the neutrino floor. For example, whenever $\rho_X \gtrsim 10^{12}$ GeV cm$^{-3}$, the identification of any 100$M_{\odot}$ would probe DM-proton cross sections below the neutrino floor. In the case of a 1000$M_{\odot}$ Pop III star, the corresponding $\rho_{X,\text{mf}} \sim 10^{16}$ GeV cm$^{-3}$.

In summary, we demonstrated that the observation of Pop III stars can be used to place strong constrains on the DM-proton cross section. Applying our method to the candidate Pop III system at $z \sim 7$ [33], we obtain some of the most stringent bounds to-date. Followup observations with JWST are necessary to confirm the Pop III nature of the system we used here, and therefore the limits we obtained.

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