

The impact of fuzzy dark matter dynamics on the accumulation and fragmentation of primordial gas

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ABSTRACT

Fuzzy dark matter (FDM), particularly in the 10^{-22} eV mass regime, is frequently used to characterize wave-like interference effects. It exhibits macroscopic wave properties, which drive distinct baryonic dynamics within collapsed haloes. Using the hydrodynamical code AREPO with the AXIREPO module and primordial chemistry, we simulate the assembly of haloes with masses $3 \times 10^8 M_\odot \leq M_h \leq 8 \times 10^9 M_\odot$ across a range of axion masses $1 \times 10^{-22} \text{ eV} \leq m_a \leq 7 \times 10^{-22} \text{ eV}$. We investigate how small-scale dynamics of the FDM density field affect the accumulation of cold, dense gas essential for primordial star formation. We demonstrate that gas collapse is suppressed by a twofold mechanism: a delay driven by the geometry of the FDM solitonic core and a secondary dynamical barrier caused by stochastic wave fluctuations. While the flattened solitonic potential profile itself inhibits central gas accumulation, these wave-driven dynamics provide a further layer of disruption and angular momentum support, which in certain regimes prevents gas from reaching the central, compact, high-density configurations characteristic of cold dark matter. Consequently, sites of star formation are shifted away from a single central peak toward a population of lower mass clusters. Our work provides a physical framework for calibrating halo mass-dependent star formation efficiencies in FDM cosmologies, where internal processes may delay Cosmic Dawn beyond the effects of the initial power spectrum cut-off. These results are essential for interpreting realistic observational constraints from future 21-cm signal observations and the faint-end luminosity functions observed by the *James Webb Space Telescope* within FDM models.

Key words: methods: numerical – stars: Population III – galaxies: haloes – galaxies: high-redshift – dark ages, reionization, first stars – dark matter.

1 INTRODUCTION

Among the candidates for dark matter (DM), the standard cold dark matter (CDM) paradigm has been remarkably successful in explaining the large-scale structure of the Universe (V. Springel et al. 2005). In the standard CDM model, typical candidates are weakly interacting massive particles (WIMPs) or quantum chromodynamic (QCD) axions, with characteristic particle masses ranging from GeV to TeV scales for thermally produced WIMPs and μeV to eV scales for axions (whose production mechanism is non-thermal). Historically, fuzzy dark matter (FDM) with an axion-like particle mass of $m_a \sim 10^{-22}$ eV was proposed as an alternative to CDM that maintained these large-scale successes while potentially resolving persistent challenges observed on sub-galactic scales (M. Boylan-Kolchin, J. S. Bullock & M. Kaplinghat

2011; D. H. Weinberg et al. 2015; A. Del Popolo, F. Pace & M. Le Delliou 2017). These discrepancies, often referred to as the ‘cusp-core’ problem, the ‘too big to fail’ problem, and the ‘missing satellites’ problem (D. H. Weinberg et al. 2015; D. J. E. Marsh 2016; J. S. Bullock & M. Boylan-Kolchin 2017; L. Hui et al. 2017; A. Eberhardt & E. G. M. Ferreira 2025), involve an overprediction of the density and abundance of small-scale structures in DMO simulations. The status of these problems has shifted in recent years, with many of these issues now largely considered to be resolvable within the CDM framework by accounting for baryonic feedback processes. For example, supernova-driven gas outflows and reionization can flatten central density profiles and suppress the formation of low-mass satellites (e.g. A. R. Wetzel et al. 2016; S. Garrison-Kimmel et al. 2019; L. V. Sales, A. Wetzel & A. Fattahi 2022). Moreover, models in which FDM comprises 100 per cent of DM have been significantly constrained by modern observations. High-resolution probes of the Lyman- α forest currently disfavour monolithic models with $m_a \lesssim 2 \times 10^{-20}$ eV

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(V. Iršič et al. 2017; E. Armengaud et al. 2017; T. Kobayashi et al. 2017; K. K. Rogers & H. V. Peiris 2021), while studies of dynamical heating in ultra-faint dwarfs have pushed these lower bounds as high as $m_a \gtrsim 3 \times 10^{-19}$ eV, or even $m_a \gtrsim 8 \times 10^{-18}$ eV if objects like Ursa Major III/Ultraviolet Near Infrared Optical Northern Survey (UNIONS) 1 are confirmed as being galaxies (N. Dalal & A. Kravtsov 2022; S. May, N. Dalal & A. Kravtsov 2025).

Despite this conceptual shift and the mounting observational evidence, it is premature to fully exclude the $m_a \sim 10^{-22}$ eV regime as a physical possibility, because the observational landscape for dwarf galaxies remains complex. While some systems suggest high m_a bounds, other recent analyses of DM-dominated dwarf populations continue to find evidence favouring the $m_a \sim 10^{-22}$ eV mass range (M. Indjin et al. 2026a; K. Korshynska et al. 2026). This suggests that the observational signatures of FDM may be highly sensitive to the specific evolutionary history of individual haloes. Moreover, FDM with $m_a \sim 10^{-22}$ eV remains a highly appealing and widely studied DM scenario owing to its unique, testable predictions arising from fundamental physics (L. Hui et al. 2017; D. Dutta Chowdhury et al. 2023; S. May & V. Springel 2023; P. Mocz et al. 2023; I.-K. Liu, N. P. Proukakis & G. Rigopoulos 2023; X. Shen et al. 2024; P.-Y. Liao et al. 2025; A. Eberhardt & E. G. M. Ferreira 2025; I. Álvarez-Ríos & F. S. Guzmán 2025; E. O. Nadler et al. 2025) as well as the limiting case of mixed dark matter (MDM) models in which FDM comprises a fraction of DM (B. Schwabe et al. 2020; A. Laguë et al. 2024; T. Dome et al. 2025; Y. F. Wang 2026). Motivated by these considerations, we simulate monolithic FDM in the $m_a = 1 \times 10^{-22}$ to 7×10^{-22} eV range to explore the interplay between wave-driven dynamics and primordial collapsing gas. Further motivation is provided in Section 2 of this paper.

In the classical monolithic FDM scenario, the DM consists of ultra-light bosonic particles, often identified as ultra-light axions, with masses on the order of 10^{-22} eV (W. Hu, R. Barkana & A. Gruzinov 2000; D. J. E. Marsh & J. Silk 2014; E. G. M. Ferreira 2021; P.-H. Chavanis 2025). These particles are similar to those derived from the axion theory initially developed to solve the strong charge-parity (CP) problem in quantum chromodynamics. From a theoretical particle physics perspective, such ultra-light axion-like particles spanning a wider mass range of 10^{-25} eV $\lesssim m_a \lesssim 10^{-18}$ eV arise naturally in top-down string theory compactifications (e.g. A. Arvanitaki et al. 2010). The extremely low mass of FDM particles results in macroscopic wave effects, where their de Broglie wavelength is on $\sim$ kpc scales, comparable to the sizes of dwarf galaxies. The wave-like nature of FDM introduces novel dynamics into the behaviour of DM on astrophysical scales: the very large de Broglie scale provides ‘quantum pressure’ that naturally smooths central cusps and suppresses the formation of (sub)haloes below a certain mass threshold. In the context of macroscopic structure formation, FDM is not treated as a discrete collection of particles, but rather modelled as a coherent, non-relativistic scalar field governed by the Schrödinger–Poisson (SP) equations (W. Hu et al. 2000; D. J. E. Marsh 2016).

On cosmological scales, the suppression of small-scale structures in FDM models results in fewer low-mass DM haloes forming compared to predictions from CDM (e.g. M. Kulkarni et al. 2022; S. May & V. Springel 2023). This has direct implications for the cooling and condensation of gas within these haloes, which are the sites of early star formation. In the standard CDM paradigm, first star (Population III) formation occurs in ‘minihaloes’ of 10^5 to $10^6 M_\odot$ where gas collapse is driven by molecular

hydrogen (H_2) cooling (V. Bromm & R. B. Larson 2004; S. Glover 2013). Because FDM suppresses these small-scale fluctuations, primordial star formation is shifted to more massive atomic cooling haloes ($M_h \gtrsim 10^7$ to $10^9 M_\odot$), where the virial temperature reaches 10 000 K. At these temperatures, the initially hot gas can cool down via efficient $\text{Ly}\alpha$ emission, initiating the runaway collapse phase ($n > 10^2 \text{ cm}^{-3}$). In this work, we employ ‘sink particles’ as a numerical proxy for this cold dense gas, representing the pre-stellar collapse phase and allowing us to track the star formation potential of each halo.

To date, there have been many studies that investigate the large-scale properties of FDM cosmologies, both with and without gas (e.g. T.-P. Woo & T. Chiueh 2009; H.-Y. Schive, T. Chiueh & T. Broadhurst 2014a; J. Veltmaat, J. C. Niemeyer & B. Schwabe 2018; P. Mocz et al. 2017, 2019, 2020; S. May & V. Springel 2021, 2023; T. Dome et al. 2023a, b, 2025; T. Dome, R. Azhar & A. Fialkov 2024). However, few studies have captured the smaller scale effects of the FDM dynamics on the collapsing gas and subsequent star formation, which we consider in this paper. While earlier works such as S. Hirano, J. M. Sullivan & V. Bromm (2018) looked at the global shift in the timing and sites of first star formation in FDM cosmologies, and M. Kulkarni et al. (2022) or Y.-M. Yang, X.-J. Bi & P.-F. Yin (2024) considered how modified potential profiles and interference patterns might inhibit gas condensation and stellar density profiles, our innovation lies in capturing the full interplay between the wave-interference dynamics, the primordial chemistry of the gas, and the gas physics of collapse. The wave-like nature of FDM, including phenomena like soliton formation and interference patterns, leads to granular density fluctuations in the DM on small scales (W. Hu et al. 2000). We might expect that these persistent dynamics would result in induced turbulence within the collapsing gas, which has been shown in the context of CDM simulations (e.g. R. S. Klessen et al. 2010; A. Stacy, V. Bromm & A. Loeb 2011; M. R. Krumholz et al. 2018; A. Soam et al. 2024) to be a key factor in star formation processes.

The gravitational effects of these fluctuations on stellar dynamics have been studied through simulations of local structures (N. C. Amorisco & A. Loeb 2018; B. V. Church, P. Mocz & J. P. Ostriker 2019; D. J. E. Marsh & J. C. Niemeyer 2019; N. Dalal & A. Kravtsov 2022; S. May et al. 2025). These studies demonstrate that collapsed FDM structures can induce observable effects on embedded baryons. However, to date, there has been little work on how these dynamics affect the collapse phase in the earliest haloes.

Calibrating these star formation models is essential for developing a crucial bookend case for the more complex MDM models as well as producing more reliable predictions for high-redshift galaxy observations with *James Webb Space Telescope* (JWST) and 21-cm experiments like the Square Kilometer Array. While JWST has made revolutionary measurements of high-redshift galaxies, particularly at the bright and massive end of the ultraviolet luminosity function (UVLF; S. L. Finkelstein et al. 2023), observations from surveys like the Hubble Frontier Fields are providing constraints on the faint end of the UVLF, which are most sensitive to FDM particle masses (J. Sippl et al. 2025). A delay or disruption in early star formation could also have observable consequences for the 21-cm signal, potentially placing lower limits on the FDM particle mass (A. Lidz & L. Hui 2018; O. Nebrin, R. Ghara & G. Mellema 2019; J. Flitter & E. D. Kovetz 2022). Importantly, these 21-cm studies typically only account for the suppression of the

halo mass function, thus ignoring the additional suppression in star formation introduced by small-scale FDM dynamics; we aim to improve on this by investigating how the intrinsic FDM dynamics further impact the small-scale star formation efficiency.

This work aims to bridge the existing gap in the literature by considering the impact of FDM dynamics on primordial star-forming gas across a range of axion-like particle masses ($m_a = 1 \times 10^{-22}$ to 7×10^{-22} eV) and DM halo masses ($M_h = 3 \times 10^8$ to $8 \times 10^9 M_\odot$). We create a series of simulated FDM haloes using the AXIREPO code within the framework of AREPO (V. Springel 2010; S. May & V. Springel 2021, 2023) with the addition of a primordial chemistry network (P. C. Clark et al. 2011), which we use to model both the atomic and molecular hydrogen (H_2) cooling necessary to track first star formation in FDM environments. A summary of our simulations can be found in Table 1. This paper is organized as follows: in Section 2, we discuss the continuous motivation for FDM; in Section 3, we outline the FDM theory, followed by methodology in Section 4. We then present our results in Section 5 and discuss the implications in Section 6. A summary of our conclusions can be found in Section 7.

2 MOTIVATION

FDM remains a highly appealing DM scenario which yields numerous unique, testable predictions arising from fundamental physics (e.g. A. Eberhardt & E. G. M. Ferreira 2025) and providing the boundary case for the increasingly prominent MDM models (e.g. A. Laguë et al. 2024; T. Dome et al. 2025).

Current dwarf galaxy observations present a diverse and sometimes contradictory picture of the axion mass. While many frameworks rule out $m_a \sim 1 \times 10^{-22}$ eV as a monolithic candidate, the continued emergence of signatures favouring this scale (M. Indjin et al. 2026a; K. Korshynska et al. 2026) points toward a more nuanced reality. It is likely that the resulting observational footprints are sensitive to the specific evolutionary pathways of individual structures, rather than being a simple function of the global DM model.

Furthermore, and most critical for this work, the 10^{-22} eV mass range serves as a fundamental benchmark for more complex, viable cosmologies such as MDM or ‘axiverse’ scenarios (A. Arvanitaki et al. 2010; M. Gosenca et al. 2023; A. Laguë et al. 2024; A. Eberhardt & E. G. M. Ferreira 2025; T. Dome et al. 2025). From a particle physics perspective, the presence of only a single ultra-light DM species, although most commonly used in studies of structure formation for reasons of simplicity and efficiency, would indeed be rather unlikely. By relaxing this assumption of a monolithic DM sector, MDM models provide a physically motivated framework to potentially reconcile the strongest large-scale structure constraints from the Lyman- α forest (K. K. Rogers & H. V. Peiris 2021) and halo mass functions (J. Sipple et al. 2025) with the stringent dynamical bounds derived from local dwarf galaxies (N. Dalal & A. Kravtsov 2022; S. May et al. 2025). As demonstrated by T. Dome et al. (2025), MDM haloes exhibit a halo mass-dependent axion depletion effect where low-mass haloes (precisely those used to derive the most stringent bounds $> 10^{-19}$ eV) are significantly depleted of their axion component relative to the cosmic mean. These systems are effectively CDM-like in their internal structure, significantly suppressing the efficiency of any dynamical heating relative to the monolithic case, while still allowing for a significant $m_a \sim 10^{-22}$ eV component to exist globally.

Studies of classical monolithic FDM in the $m_a = 1 \times 10^{-22}$ to 7×10^{-22} eV range are further motivated by scalability. The FDM dynamics are governed by the SP equations, which admit a well-known scaling relation (see e.g. P. Mocz et al. 2017). Under this scaling, the dynamics of a well-understood halo can be mapped on to a new system with a different halo mass and m_a . While it is important to note that baryonic physics (such as gas cooling or stellar feedback) do not scale in this way and require independent investigation, the gravitational behaviour of the FDM remains qualitatively similar across scales. Therefore, resolving the regime where the dynamics are most prominent is a necessary prerequisite to understanding the effects that FDM dynamics can have on the primordial gas. By isolating these mechanisms at computationally resolvable systems such as explored in this work, we can establish a pathway towards for the broader class of wave-DM models at different m_a , regardless of whether the axion is a monolithic or subdominant DM component.

Considering this strong continuous motivation, we simulate ‘classical monolithic’ FDM models in the $m_a = 1 \times 10^{-22}$ to 7×10^{-22} eV range in a controlled and computationally feasible setup to characterize the interactions between the FDM dynamics and the baryonic sector. We explore the implications of the wave-driven dynamics on the collapse and distribution of cold, potentially star-forming, gas. With this work, we aim to provide the community with a crucial bookend case for the more complex MDM models and to enable new independent FDM/MDM constraints at Cosmic Dawn.

3 THEORETICAL BACKGROUND

For this work, we consider a physical system composed of two primary components: the DM, which provides the underlying gravitational potential and cosmic scaffolding, and the primordial gas, which accretes into these potential wells. The physical problem centres on how the unique small-scale dynamics of FDM modify the standard process of gas cooling and collapse. While the DM serves as the gravitational driver, its wave-like nature in the FDM model creates a non-static potential, characterized by a central solitonic core and persistent interference-induced fluctuations. Our setup investigates the ‘active’ role of these FDM dynamics, focusing on how they exert time-varying gravitational influences that can hinder gas accretion and alter the timing of the first star formation events in the early Universe.

In the following section, we lay out the fundamental theory underlying the FDM model (Section 3.1), and the important dynamics within FDM haloes (Section 3.2), both of which are key in understanding the resulting effects on gas collapse.

3.1 Ultra-light dark matter

In the non-relativistic limit and at zero temperature, a spin-0 ultra-light scalar field represented by the (collective, mean-field) ‘wavefunction’ ψ is governed by the SP equations (W. Hu et al. 2000):

$$i\hbar \frac{\partial \psi}{\partial t} = -\frac{\hbar^2}{2m_a} \nabla^2 \psi + m_a V \psi, \quad (1)$$

$$\nabla^2 V = 4\pi G (\rho - \bar{\rho}). \quad (2)$$

In the above equations, ψ is the FDM ‘wavefunction’ normalized such that the DM density $\rho_{\text{dm}} = |\psi|^2$, m_a is the boson mass, and V is the gravitational potential which is affected by FDM as well as the baryonic component via the Poisson equation (2), where

Table 1. Summary of simulations run. FDM runs are characterized by a fixed spatial grid resolution $\Delta x = 31.25$ pc (in physical pc), while CDM runs are characterized by a mass resolution of 480, 180, 48, and $18 M_\odot$ for our largest to smallest haloes (following the first row of the table). The dashes ‘–’ represent cases that were omitted (e.g. due to tidal instability or being outside the relevant parameter space).

Halo mass ($M_\odot$)	FDM models (m_a)								CDM
	1×10^{-22} eV		2×10^{-22} eV		3×10^{-22} eV		7×10^{-22} eV		
	Dynamic	Frozen	Dynamic	Frozen	Dynamic	Frozen	Dynamic	Frozen	
8×10^9	✓	✓	–	–	✓	✓	✓	✓	✓
3×10^9	✓	✓	✓	✓	✓	✓	✓	✓	✓
8×10^8	–	–	✓	✓	✓	✓	✓	✓	✓
3×10^8	–	–	–	–	✓	✓	✓	✓	✓

the total matter density is given by $\rho = \rho_{\text{dm}} + \rho_{\text{gas}}$ and $\bar{\rho}$ is the mean matter density.

Instead of the wavefunction, the axion-like field can be represented by a fluid with density and velocity, related to the complex wavefunction via the Madelung transformation (E. Madelung 1927; A. Suárez & T. Matos 2011; P.-H. Chavanis & L. Delfini 2011), where we decompose ψ into an amplitude $\sqrt{\rho_{\text{dm}}}$ and a phase S . Given this decomposition, we now have

$$\begin{aligned}\psi &= \sqrt{\rho_{\text{dm}}} e^{iS/\hbar}, \\ v &= \nabla S / m_a.\end{aligned}\quad (3)$$

The SP equations obey both conservation of mass:

$$M = \int \rho_{\text{dm}} d^3x, \quad (4)$$

and energy:

$$\begin{aligned}E_{\text{tot}} &= \int \left[\frac{\hbar^2}{2m^2} |\nabla \psi|^2 + \frac{1}{2} V |\psi|^2 \right] d^3x \\ &= \int \frac{\hbar^2}{2m^2} (\nabla \sqrt{\rho_{\text{dm}}})^2 d^3x + \int \frac{1}{2} \rho_{\text{dm}} v^2 d^3x + \int \frac{1}{2} \rho_{\text{dm}} V d^3x \\ &= K_\rho + K_v + W,\end{aligned}\quad (5)$$

where the integration volume is any closed volume. The total kinetic energy $K = K_\rho + K_v$ is given by the sum of the energy due to the quantum pressure tensor K_ρ and the classical contribution K_v . W is the potential energy. The decomposition of the kinetic energy term is obtained by substituting the Madelung form into the gradient $|\nabla \psi|^2$. The phase information S is mapped directly on to the bulk velocity term K_v via its gradient (∇S), while the absolute value of the wavefunction amplitude yields the quantum pressure term K_ρ . Note that in the potential energy term W , the phase factor $e^{iS/\hbar}$ disappears entirely when taking the inner product $|\psi|^2 = \psi \psi^*$, leaving only the dependence on the density ρ_{dm} .

For m_a below approximately 10^{-18} eV, the FDM density field exhibits wave behaviour on astrophysical scales. The characteristic scale of these fluctuations is the de Broglie wavelength, $\lambda_{\text{dB}} = \hbar / (m_a \sigma)$. Here, σ represents the effective velocity dispersion of the system. It is important to note that because the Madelung velocity v (equation 3) is derived from a single-valued phase S , the local velocity dispersion is zero by definition. In this context, σ is understood as a macroscopic quantity representing the root-mean-square (RMS) velocity of the wave-field fluctuations, or equivalently, as a measure of the total kinetic energy within a given volume. For a typical velocity dispersion of $\sigma = 100 \text{ km s}^{-1}$ and $m_a = 10^{-22}$ eV, this yields a de Broglie scale (in physical

units) of:

$$\lambda_{\text{dB}} \approx 4 \left(\frac{10^{-22} \text{ eV}}{m_a} \right) \left(\frac{100 \text{ km s}^{-1}}{\sigma} \right) \text{ kpc}. \quad (6)$$

This wave-like behaviour results in distinctive observational signatures, most notably the suppression of small-scale structure formation (W. Hu et al. 2000; L. Hui et al. 2017). Within the centres of FDM haloes, quantum pressure counteracts gravitational collapse and introduces a ‘quantum Jeans scale’ (M. I. Khlopov, B. A. Malomed & I. B. Zeldovich 1985) which leads to the formation of solitonic cores (H.-Y. Schive et al. 2014a; P. Mocz et al. 2017). These cores represent stable, high-density central regions where the density profile is maintained by the balance between gravity and quantum pressure, and can be approximated as:

$$\rho_{\text{dm}}(r) = \frac{\rho_c}{\left(1 + 0.091 \left(\frac{r}{r_c} \right)^2 \right)^8}, \quad (7)$$

where ρ_c is the central density and r_c is the core radius. The factor of 0.091 adjusts the profile such that $\rho_{\text{dm}}(r_c) \approx 0.5 \rho_c$. This core and its properties dominate the central potential of an FDM halo and, therefore, regulate the collapse of the gas in its central regions.

3.2 Dynamics of FDM haloes

One of the unique signatures of the FDM physics is that the wavefunction interference produces ubiquitous oscillations in the local density field across the halo. These oscillations have an amplitude of order unity $\delta \rho_{\text{dm}} \sim \rho_{\text{dm}}$ and characteristic length- and time-scales of $r \approx \lambda_{\text{dB}} / 2\pi = \hbar / m_a \sigma$ and $\tau_{\text{osc}} \approx \lambda_{\text{dB}} / v = \hbar / m_a \sigma^2$ where λ_{dB} is as defined above in Section 3.1. For a typical combination of $m_a = 10^{-22}$ eV and $\sigma = 100 \text{ km s}^{-1}$ from our work, this results in:

$$r \approx 0.6 \left(\frac{10^{-22} \text{ eV}}{m_a} \right) \left(\frac{100 \text{ km s}^{-1}}{\sigma} \right) \text{ kpc}, \quad (8)$$

$$\tau_{\text{osc}} \approx 40 \left(\frac{10^{-22} \text{ eV}}{m_a} \right) \left(\frac{100 \text{ km s}^{-1}}{\sigma} \right)^2 \text{ Myr}. \quad (9)$$

This phenomenon results in corresponding fluctuations in the gravitational potential within the halo (B. V. Church et al. 2019; D. J. E. Marsh & J. C. Niemeyer 2019; I. Álvarez-Ríos & F. S. Guzmán 2023).

The soliton at the centre of each FDM halo is also not static, but instead exhibits similar order-unity fluctuations in its central density (with a period of $\sim \tau_{\text{osc}}$) as well as a random walk behaviour about the centre of mass of the halo with magnitude $\sim \lambda_{\text{dB}}$ (H.-Y. Schive et al. 2014a, 2020; L. Hui 2021). These behaviours are a result of persistent interference of the wavefunction outside

the soliton with the soliton itself. This is important as the soliton dominates the gravitational potential at small distances from the halo centre. We expect these effects are most pronounced in small haloes in models with a low m_a , since the mass contained within the soliton scales roughly as $M_s \propto M_h^{1/3} m_a^{-1}$ (H.-Y. Schive et al. 2014b, although this relation shows significant variation, see e.g. H. Y. J. Chan et al. 2022) and thus the proportion of halo mass contained in the soliton can approximately be estimated as $M_s/M_h \propto M_h^{-2/3} m_a^{-1}$.

To illustrate these theoretical principles, we first examine a simplified dark matter-only (DMO) configuration. This allows us to isolate the dynamical behaviour of DM and confirm the theory results we have discussed before introducing the complexities of gas physics. In this simplified setup, the gravitational potential is determined solely by the FDM density, allowing us to isolate the wave-driven fluctuations of the core without the baryons which will be discussed in Section 5. For this illustrative case, we consider an isolated halo of mass $M_h \approx 3 \times 10^9 M_\odot$ with an axion mass $m_a = 10^{-22}$ eV within a 32 kpc cubic box and a DM resolution of $1/32 \text{ kpc} = 31.25 \text{ pc}$.

The evolution of the soliton can clearly be seen in these simulations. Fig. 1(a) shows a series of snapshots from a DMO halo zoomed-in on the central region, where the green line shows the movement of the centre of the solitonic core sampled approximately every 1 Myr over the course of four snapshots separated by approximately 100 Myr. Each zoomed panel shows the same simulation box coordinates for consistency between the panels. The panel marked (b) shows how the central density ρ_c , defined as the density of the densest point of the DM, evolves over time, with order unity fluctuations in the central density contrast $(\rho_c - \langle \rho_c \rangle) / \langle \rho_c \rangle$ clearly visible, where $\langle \rho_c \rangle$ is the time-averaged central density of the halo. The oscillation time-scale τ_{osc} closely matches the theory value (shown as the dashed vertical line) and for this halo with $m_a = 1 \times 10^{-22}$ eV and $\sigma \approx 80 \text{ km s}^{-1}$ corresponds to $\tau_{\text{osc}} \approx 50 \text{ Myr}$.

In this work, the centre of each DM halo is defined as the minimum of its gravitational potential unless otherwise specified. This approach is adopted in preference to its the centre of mass because it is a more physically relevant quantity that governs the dynamics of the central gas. Since the potential fluctuations at the centre of the halo are the primary mechanism we hypothesize to affect gas collapse and star formation, using the centre of potential provides a physically motivated and direct reference frame for analysing the response of the gas to these small-scale dynamics. In FDM, this point undergoes a stochastic random walk due to soliton migration, as shown in Fig. 1(a). All radial profiles and analyses presented in this paper are centred on this time-varying point unless explicitly stated otherwise.

4 METHODOLOGY

In this section, we describe the numerical framework and experimental design used to investigate gas collapse processes in the first FDM haloes. Our methodology is structured to isolate the physical effects of wave-driven fluctuations from standard gravitational collapse.

First, in Section 4.1, we outline our simulation design, including the three-way comparison between CDM, FDM, and ‘frozen’ FDM models. The computational details of the AREPO and AXIREPO solvers, including the gravitational coupling between DM and baryons, are detailed in Section 4.2. Lastly, we describe the generation of our initial conditions and the physical

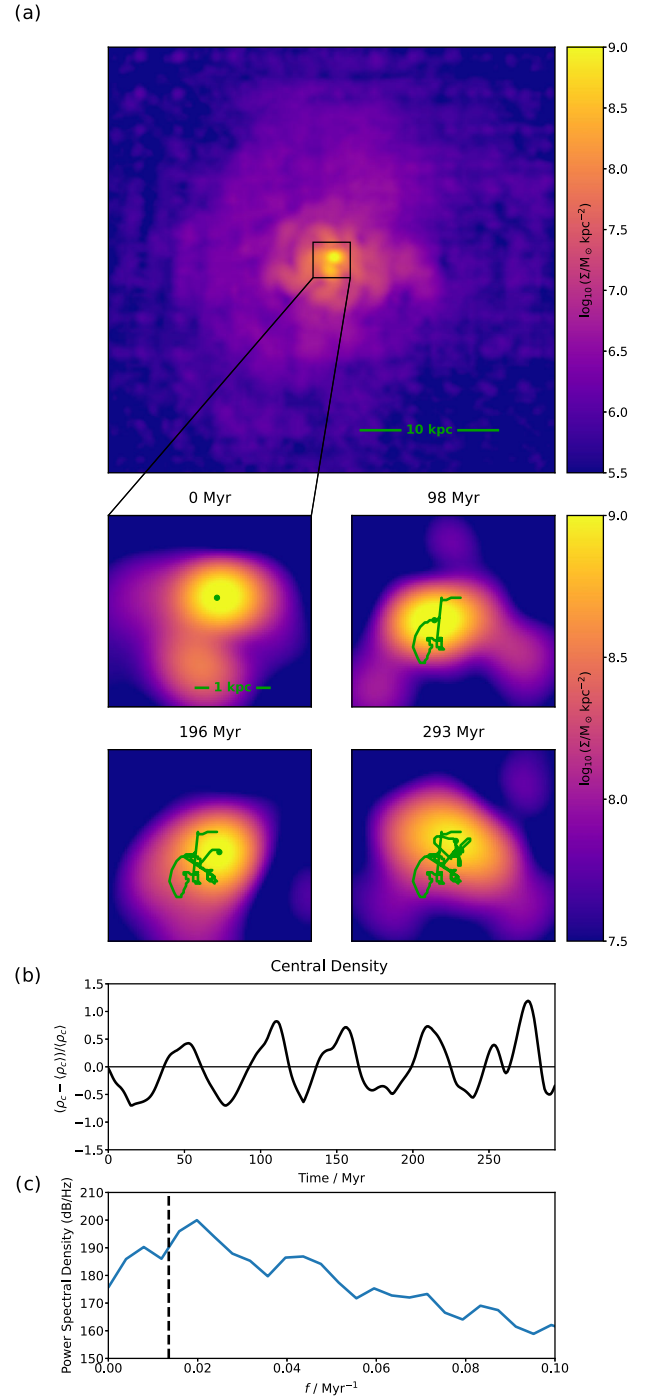


Figure 1. Dynamics of the solitonic core at the centre of a simulated FDM halo with a halo mass $M_h = 3 \times 10^9 M_\odot$ for axion mass $m_a = 10^{-22}$ eV. The panels marked (a) show the projected (log) density of the DM ρ_{dm} of a zoomed-in region at the centre of an FDM halo for a series of snapshots separated by ~ 100 Myr. The line shows a track of the central density from its initial position at 0 Myr within the halo sampled with time resolution $\Delta t \approx 1$ Myr at four snapshots throughout the simulation, spaced approximately 100 Myr apart. The panel marked (b) shows the order unity fluctuations of the central density along this track over time, with a characteristic period of $f \sim 0.02 \text{ Myr}^{-1} \Rightarrow \tau_{\text{osc}} \sim 50 \text{ Myr}$, which can be identified by the peak in the power spectral density of this oscillation in the bottom panel marked (c). This lines up with the theoretical value for the period of oscillation $\tau_{\text{osc}} \approx 65 \text{ Myr}$ at the centre of the halo, shown by the dashed line.

motivation for our choice of halo and axion-like particle masses in Section 4.3.

4.1 Simulation design and objectives

The primary goal of this work is to isolate and quantify how the unique small-scale dynamics of FDM, specifically the fluctuations of the central solitonic core, influence the collapse of primordial gas and the assembly of the dense star-forming regions. To achieve this, we compare the evolution of gas infall on to an isolated DM halo across three distinct DM scenarios: standard CDM, fully simulated ‘dynamic’ FDM, and a ‘frozen’ FDM model (explained below). For the remainder of this work ‘dynamic’ FDM refers to those runs where the SP equations (equations 1 and 2) are solved and the FDM evolves over time.

For our simulation suite, summarized in Table 1, we chose four halo masses at 3×10^8 , 8×10^8 , 3×10^9 , and $8 \times 10^9 M_\odot$. We then evolved each of these initial setups for four axion masses (10^{-22} , 2×10^{-22} , 3×10^{-22} , and 7×10^{-22} eV), as well as the CDM comparison case. All simulations are performed in fixed physical boxes with side length $L_{\text{box}} = 32$ kpc and with periodic boundary conditions which are necessary for the pseudo-spectral SP solver (see Section 4.2.1 and P. Mocz et al. 2017; S. May & V. Springel 2021). A box size of $L_{\text{box}} = 32$ kpc is sufficient for our needs; at redshifts $z \sim 10$ to 20, the virial radii for the haloes in our sample range from a few hundred pc to a few kpc (G. L. Bryan & M. L. Norman 1998; R. Barkana & A. Loeb 2001). Thus, our 32 kpc box means that the isolated system remains well-defined. This choice is also motivated by our need to resolve the internal dynamics of the central $\sim$ kiloparsec. The DM resolution is handled according to the specific solver:

(i) *FDM resolution.* All FDM runs (both ‘dynamic’ and ‘frozen’) utilize a fixed 1024^3 Cartesian grid, providing a uniform spatial resolution of $\Delta x_{\text{dm}} = 31.25$ pc. Numerical convergence in FDM simulations requires that the solitonic core radius be resolved by at least four grid cells to accurately capture the dynamics of the soliton and central density fluctuations (P. Mocz et al. 2017; S. May & V. Springel 2021).

(ii) *CDM resolution.* Our CDM simulations use 256^3 N -body particles. This corresponds to a mass resolution (m_{DM}) ranging from $\sim 18 M_\odot$ for our $3 \times 10^8 M_\odot$ halo to $\sim 480 M_\odot$ for our $8 \times 10^9 M_\odot$ halo. To resolve the central density profile, we employ a gravitational softening length of $\epsilon = 3.1$ pc. This scale is an order of magnitude smaller than the fixed spatial grid resolution used in our FDM simulations, ensuring that the central cusp is accurately captured without numerical flattening. This disparity in resolution is intentional, as it allows the CDM cases to provide a true cusp at the halo centre, whereas the flattened FDM density profile does not require such high resolution.

(iii) *Gas resolution.* The gas component is initialized with a total mass in line with the cosmological baryon fraction and using 256^3 gas cells. Initially this corresponds to a gas cell mass of $m_{\text{gas}} \approx 3 M_\odot$ for our lowest halo mass, and up to $m_{\text{gas}} \approx 90 M_\odot$ for our highest halo mass. To accurately capture the fragmentation of primordial gas and avoid numerical artifacts, we employ a Jeans refinement criterion. We ensure that the local Jeans length is resolved by a minimum of 16 cell diameters throughout the simulation. This exceeds the standard Truelove criterion (J. K. Truelove et al. 1997) by a factor of four, ensuring that the fragmentation observed in the solitonic cores is physically driven by the interplay of gravity, thermal pressure, and FDM-induced fluctuations.

Each halo was generated through the merging of a number of solitons within the simulation box, which are allowed to relax and virialize to form one halo (further details in Section 4.3.1). For consistency and ease of comparison between models, the soliton merger initial conditions for a given halo mass across all three models (FDM, ‘frozen’ FDM, and CDM) are the same.

To ensure these results are numerically robust, we performed convergence tests on a 2048^3 grid for the $M_h = 8 \times 10^9 M_\odot$ case. Even with the highest axion mass ($m_a = 7 \times 10^{-22}$ eV), which presents the smallest physical substructures to resolve, we saw < 2 per cent deviation in key metrics used in this paper between the 1024^3 and 2048^3 grids.

The choice to use a box of fixed physical size was motivated by two factors:

(i) In a box of fixed physical size, our fixed-grid FDM spatial resolution is not degraded over time, as the universe expands.

(ii) For studying individual haloes, once they have collapsed, they have become decoupled from the cosmological expansion and so the expansion should not affect the internal dynamics of the halo.

The wave-mechanical nature of FDM imposes a strict constraint on the simulation time-step; to evolve the phase of the scalar field correctly, the maximum allowed time-step satisfies $\Delta t \propto \Delta x^2$ (the Schrödinger Courant condition), which is significantly more restrictive than the standard gravitational Courant condition used in our CDM runs (P. Mocz et al. 2017; S. May & V. Springel 2021). Resolving these dynamics while simultaneously tracking gas fragmentation at high precision would be computationally prohibitive in a larger cosmological box.

The ‘frozen’ FDM model serves as a diagnostic tool; in this setup, we first evolve a DMO FDM halo until it reaches a virialized state. We then pause the DM part of the simulation, thereby ‘freezing’ it in place. This ‘frozen’ state is then used to generate a static background gravitational potential. In these simulations, the gas is allowed to cool and collapse within this static cored potential. Since the DM potential is fixed, the gas does not experience the order-unity density fluctuations or the physical migration of the soliton centre obtained using the full-physics SP solver. We expect that this ‘frozen’ potential will lead to more efficient gas accumulation than the ‘dynamic’ FDM case. However, we still expect to see suppression in gas infall and star formation relative to CDM due to the presence of the solitonic core.

This three-way comparison between CDM, ‘dynamic’ FDM, and ‘frozen’ FDM allows us to distinguish between effects caused by the spatial shape of the potential well (solitonic cored profile in FDM versus cusped in CDM) and those caused by the dynamic wave-like oscillations and migrations of the solitonic core and surrounding interference granules.

The range of halo masses $M_h \approx 10^8$ to $10^{10} M_\odot$ is chosen because it represents typical first star-forming haloes in FDM cosmologies; with this range of m_a , the suppression of small-scale power means that star formation is delayed until haloes reach these mass scales. These choices are motivated by both physical relevance and numerical feasibility. Previous large-scale cosmological simulations of FDM (e.g. H.-Y. Schive et al. 2014a; S. May & V. Springel 2023; T. Dome et al. 2023b) indicate that this halo mass range represents the bulk of first star formation in FDM cosmologies with $m_a \sim 10^{-22}$ eV. For $M_h < 10^8 M_\odot$, the FDM power spectrum cut-off typically prevents much smaller structures from forming for the given range of m_a . We also aim to examine M_h values where the resulting solitonic core constitutes a significant

fraction of the halo mass and so we use the approximate core-halo mass relation of H.-Y. Schive et al. (2014b) described in Section 3.2 to select the range of halo masses. We expect that the soliton mass fraction would correlate with the extent of any disruption to gas collapse as it dominates the central potential. This ensures that the wave-driven dynamics are a dominant gravitational force in the region where gas fragmentation and collapse occurs.

As well as computational constraints of resolving the smaller wave patterns with higher axion masses (e.g. $m_a \sim 10^{-20}$ eV), there is a diminishing cosmological motivation for exploring much higher m_a in the context of resolving small-scale discrepancies. At $m_a \gtrsim 10^{-20}$ eV, the FDM power spectrum and halo internal structures begin to converge toward the standard CDM limit (W. Hu et al. 2000; H.-Y. Schive et al. 2014a; L. Hui et al. 2017; N. Dalal & A. Kravtsov 2022; D. M. Powell et al. 2023), effectively removing the features that make ‘plain’ FDM a distinct and appealing alternative to CDM at these mass scales. Consequently, focusing on $m_a \sim 10^{-22}$ eV represents the most relevant region of the parameter space to explore. Our preliminary tests also showed that at 10^{-21} eV, the soliton’s fluctuations are sufficiently small and high frequencies that their impact on the gas is negligible when compared to a reference CDM case within our halo mass range, which further motivates this choice of m_a range as where the most significant departures from CDM are found.

By focusing on isolated systems, we can precisely control the halo mass M_h and axion mass m_a parameters, and binding energy of our haloes. We attempted to simulate the m_a and M_h combinations in Table 1 that show ‘–’, but due to the small periodic box and large core radii, these haloes were unstable and so did not form a bound virialized halo during halo generation. This can be seen as a result of the core-halo mass relation (H.-Y. Schive et al. 2014b; H. Y. J. Chan et al. 2022) described in Section 3.2, $M_s/M_h \propto M_h^{-2/3} m_a^{-1}$, which implies that for a given m_a and sufficiently low M_h , the solitonic core becomes so large and diffuse that $M_s \gtrsim M_h$ and thus forming a bound halo is impossible. In our simulations this occurred for our $M_h = 8 \times 10^8 M_\odot$ halo and $m_a \leq 1 \times 10^{-22}$ eV and for our $M_h = 3 \times 10^8 M_\odot$ halo for $m_a \leq 2 \times 10^{-22}$ eV.

Further details of our simulations, including our computational setup and our initial conditions, can be found in Sections 4.2 and 4.3, respectively.

4.2 Computational setup

For all our simulations, we use the well-known and tested hydrodynamics code AREPO (V. Springel 2010), which is well suited for resolving high-dynamic-range structures like primordial gas clouds. To model the different DM scenarios, we employ two distinct DM gravity solvers within the code.

4.2.1 Fuzzy dark matter

For the FDM simulation, we use the AXIREPO module to simulate the FDM physics (S. May & V. Springel 2021, 2023). AXIREPO makes use of a second-order pseudo-spectral ‘kick-drift-kick’ method to solve the SP equations (T.-P. Woo & T. Chiueh 2009). For this method, the wavefunction ψ and potential V are discretized on to a uniform N^3 Cartesian grid, which allows us to use the fast Fourier transform (fft) and its inverse (ifft). The potential can be calculated directly from the wavefunction ψ as $\rho_{\text{dm}} = |\psi|^2$

and $\rho = \rho_{\text{dm}} + \rho_{\text{gas}}$ as follows:

$$V = \text{ifft} \left[-\text{fft} [4\pi G (\rho - \bar{\rho})] / k^2 \right], \quad (10)$$

where k is the wavenumber for the given grid locations. The wavefunction is then evolved through a ‘kick’ of half a time-step due to the potential:

$$\psi \leftarrow \exp \left[-i \left(\frac{\Delta t}{2} \right) \left(\frac{m}{\hbar} \right) V \right] \psi. \quad (11)$$

This is followed by a full ‘drift’ step:

$$\hat{\psi} \leftarrow \text{fft} [\psi], \quad (12)$$

$$\hat{\psi} \leftarrow \exp \left[-i \Delta t \left(\frac{\hbar}{m} \right) \left(\frac{k^2}{2} \right) \right] \hat{\psi}, \quad (13)$$

$$\psi \leftarrow \text{ifft} [\hat{\psi}]. \quad (14)$$

The algorithm is then completed by performing another ‘kick’ step (equation 11) in order to finish the time evolution of one time-step Δt .

We also ran a set of simulations where we took the same FDM initial conditions, but ‘froze’ the FDM. i.e. paused the SP solver described above and stopped evolving the wavefunction ψ but preserved the gravitational contribution from the DM when the gas was added and evolved. This approach allows us to directly see whether any effects on the embedded gas are a result of the lack of a gravitational cusp in FDM haloes, and how much is caused by the continuous, time-varying gravitational fluctuations of the FDM.

4.2.2 Cold dark matter

Our CDM simulations treat DM as a collection of 256^3 collisionless ‘Monte Carlo’ particles sampling phase space. Gravitational forces are calculated using the standard TREEPM algorithm in AREPO. In this approach, long-range forces are computed on a Fourier grid according to equation (10), while short-range forces are handled via a hierarchical octree structure. This dual approach ensures that we capture the small-scale structure of CDM haloes (which lack the quantum pressure of FDM) with high precision (J. S. Bagla 2002; V. Springel et al. 2005).

To avoid numerical divergences during close-range particle encounters and to accurately capture the characteristic central ‘cusp’ of the CDM potential, we employ a gravitational softening length of $\epsilon = 3.1$ pc. We intentionally set this value to be smaller than the accretion radius of our sink particles (see Section 4.2.4). This ensures that the gravitational potential is resolved at scales below the star formation threshold, providing a physically robust potential down to the length and density scales of our sink formation without being artificially smoothed by the softening at the scales of interest.

4.2.3 Gas dynamics and gravitational coupling

While the DM is evolved on a Cartesian grid (FDM) or as ‘Monte Carlo’ particles (CDM), the gas component is modelled using the quasi-Lagrangian moving mesh approach of AREPO (V. Springel 2010; R. Weinberger, V. Springel & R. Pakmor 2020).

The gas is discretized using a Voronoi tessellation of the simulation volume, where the mesh-generating points move with the local fluid velocity. This allows the simulation to automatically increase resolution in high-density regions, such as the collapsing gas clouds at the centre of the DM haloes, without the numerical diffusion typical of fixed-grid Eulerian methods.

The gravitational coupling between the DM and gas is handled based on the respective DM solver:

(i) *FDM (frozen and dynamic)*. The gravitational potential V is calculated on the N^3 Cartesian grid using the pseudo-spectral solver as described in Section 4.2.1. At each synchronization step, this potential is interpolated from the grid on to the centre of mass of each Voronoi gas cell using a standard trilinear interpolation scheme, in the same way as in a PM solver (such as the PM part of the TREEPM method). The gradient of this interpolated potential then provides the gravitational acceleration for the gas due to the underlying DM. The rest of the gravitational contribution from the gas is done as in the CDM case (using the TREEPM algorithm, see below).

(ii) *CDM*. The gravitational acceleration for each gas cell is calculated directly by summing the contributions from the DM particles using the TREEPM algorithm (V. Springel 2010). The tree provides the short-range forces from nearby particles, while the Fourier grid handles the long-range component, ensuring that the gas ‘sees’ the same high-resolution potential, including the central cusp, as the CDM particles.

Alongside the gravitational considerations, we also employ the sink particle prescription as detailed by M. R. Bate, I. A. Bonnell & N. M. Price (1995), C. Federrath et al. (2010), R. G. Trefß et al. (2020), K. M. J. Wollenberg et al. (2020, further details in Section 4.2.4), and the primordial chemistry network of P. C. Clark et al. (2011, details in Section 4.2.5). We set the softening scale of the gas at $1.9 \times 10^{-1} \text{ pc} = r_{\text{sink}}/16$ which is the minimum size a gas cell should be able to reach based off the sink creation radius (Section 4.2.4) and Jeans refinement criteria requiring 16 cells per Jeans length.

4.2.4 Sink prescription

In order to limit the indefinite refinement of the gas cells and to introduce an easy measure of the quantity and distribution of cold dense gas we use a sink particle prescription within AREPO. In our simulations, these sink particles serve as a proxy for star-forming regions within the halo where the gas has met the requirements to continue cooling and collapse as outlined below (A. T. P. Schauer et al. 2021). The sink particle implementation we use here is the same as that introduced in R. G. Trefß et al. (2020, 2021) and K. M. J. Wollenberg et al. (2020) based off C. Federrath et al. (2010). These sink particles are formed when a gas cell that is at a local minimum of the potential satisfies the following criteria:

- (i) the gas cell has a number density $n > n_{\text{sink}} = 10^2 \text{ cm}^{-3}$;
- (ii) the local velocity field is convergent;
- (iii) the local acceleration is convergent;
- (iv) the total energy of the gas within the sink accretion radius (r_{sink}) is negative, i.e. the gas is gravitationally bound;
- (v) there are no other sink particles within r_{sink} .

These sink particles act as point masses within the simulation which are able to accrete mass from any gas cells around them that are within the sink particle’s accretion radius r_{sink} . We set the value of r_{sink} to be comparable to the Jeans length at the sink formation density threshold n_{sink} . In our simulations, gas at this density typically has a temperature of $\sim 10^2 \text{ K}$, resulting in a Jeans length of around 3 pc, and so we set $r_{\text{sink}} = 3 \text{ pc}$. Gas within a distance r_{sink} of an existing sink is accreted if it meets

the following requirements: it must be gravitationally bound to the sink and must have a density above n_{sink} .

Our choice of value for n_{sink} is motivated by our desire to effectively capture the bulk properties of the active, star-forming region within the halo centre while at the same time ensuring that our simulations remain computationally efficient.

It is important to clarify that at this density, the sinks do not represent individual Population III stars. Instead, they serve as numerical proxies for the onset of the runaway cooling phase where the free-fall time (t_{ff}) and cooling times (t_{cool}) of the gas are much less than characteristic time-scale of any of the DM dynamics τ_{osc} . These particles accurately capture the relative integrated star formation rate of the halo, and can be used to compare the accumulation of dense, star-forming gas between haloes, even if they do not resolve individual stars (N. Y. Gnedin, K. Tassis & A. V. Kravtsov 2009; M. Jeon et al. 2012; A. T. P. Schauer et al. 2021). By using these sink particles, we intentionally circumvent the need to model the full fragmentation and collapse into individual stellar objects, a process that would otherwise necessitate resolving significantly smaller scales and much higher densities, and hence would have much greater computational requirements.

4.2.5 Chemical model

In metal-free primordial gas, radiative cooling is governed by a small number of atomic and molecular processes, with the dominant cooling channel depending primarily on halo mass and virial temperature. In ‘minihaloes’ ($M_{\text{h}} \sim 10^5$ to $10^7 M_{\odot}$, $T_{\text{vir}} \lesssim 10^4 \text{ K}$), cooling is enabled by molecular hydrogen (H_2), which forms through gas-phase reactions and cools the gas via rotational–vibrational transitions, allowing collapse to temperatures of $\sim 200 \text{ K}$ (T. Abel, G. L. Bryan & M. L. Norman 2002; V. Bromm & R. B. Larson 2004; S. C. O. Glover & T. Abel 2008; V. Bromm & N. Yoshida 2011). In more massive haloes like the ones we aim to study ($M_{\text{h}} \gtrsim 10^7$ to $10^8 M_{\odot}$ at $z \sim 10$ –20), the virial temperature exceeds $\sim 10^4 \text{ K}$ and cooling is instead dominated initially by collisional excitation of atomic hydrogen and subsequent Ly α emission (S. P. Oh & Z. Haiman 2002; J. H. Wise & T. Abel 2007; V. Bromm & N. Yoshida 2011). In these atomic-cooling haloes, atomic hydrogen cooling regulates the onset of collapse and efficiently condenses the gas until temperatures drop below $\sim 8000 \text{ K}$. Below this threshold, atomic cooling becomes inefficient, and molecular hydrogen (H_2) cooling takes over at higher densities to guide the subsequent thermodynamic evolution and fragmentation of the gas (S. C. O. Glover & T. Abel 2008; T. H. Greif et al. 2012). Both atomic and molecular cooling are thus important regulators of cooling and collapse within our simulated systems.

The primary motivation for our choice of chemical model is to accurately track the transition from virialized gas to the first star-forming regions. We adopted a primordial chemical network and cooling function based on the ones introduced in P. C. Clark et al. (2011), with updates as described in A. T. P. Schauer et al. (2017). The model handles the primordial gas chemistry of 12 species (H , H^+ , H^- , H_2^+ , H_2 , He , He^+ , He^{++} , D , D^+ , HD , and free electrons) via 45 chemical reactions.

While AREPO handles the standard hydrodynamics, this chemistry module is essential for Population III studies as it solves the non-equilibrium rate equations for each species alongside the hydrodynamical evolution. The primary cooling channels in our simulations are driven by H_2 and HD molecules. These represent the dominant cooling mechanisms for primordial gas at

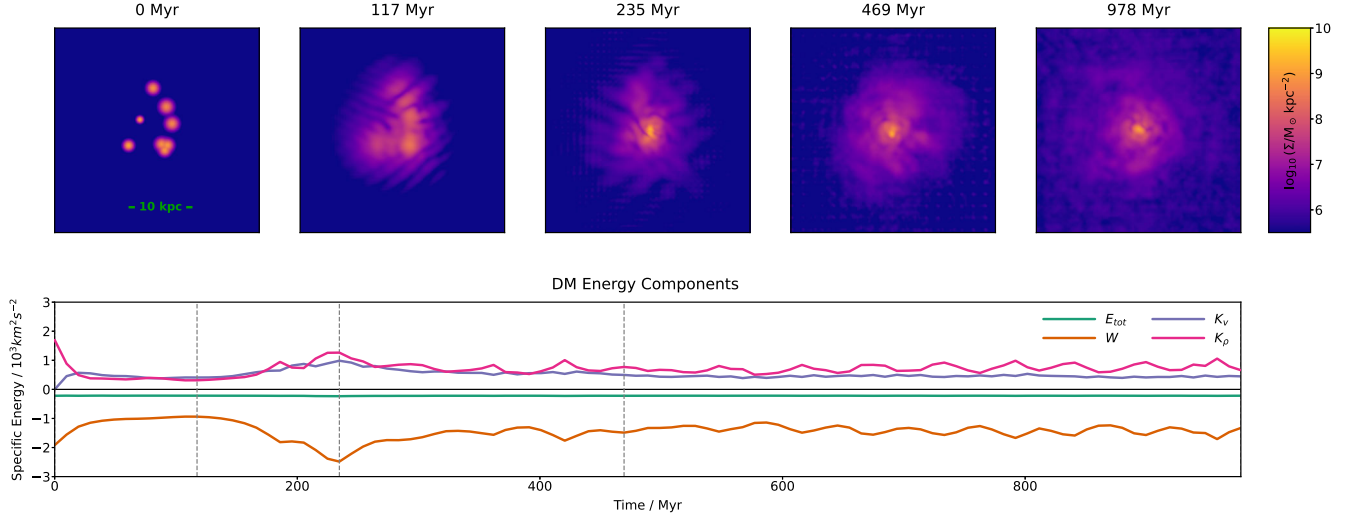


Figure 2. Top: projected (log) density snapshots of the halo assembly (DMO) starting from the initial conditions (the leftmost panel), followed by three intermediate steps at 117, 235, and 469 Myr, taken from our $3 \times 10^9 M_\odot$ halo and with $m_a = 1 \times 10^{-22} \text{ eV}$. The final halo used for our simulations is shown in the rightmost panel. Bottom: the total energy of the system (E_{tot}) and energy components: potential energy (W), classical kinetic energy (K_v) and quantum kinetic energy (K_p) throughout the assembly timeline of the halo. The vertical dashed lines correspond to the three instances shown in the intermediate top panels. The total energy is clearly conserved throughout the simulation box. We highlight the horizontal axis at zero to show clearly that the total energy of our haloes is less than zero, i.e. the haloes are bound.

temperatures below 5000–8000 K, where atomic cooling becomes inefficient (S. C. O. Glover & T. Abel 2008). In particular, HD cooling is critical for allowing the gas to surpass the H_2 cooling limit and reach temperatures below 200 K (D. Galli & F. Palla 1998), a regime that is vital for determining the fragmentation scale of the gas. The network further accounts for ionization and recombination, three-body H_2 formation, and the thermal effects of chemical composition changes, shocks, and adiabatic compression or expansion.

At each time-step, the chemical abundances are updated based on the local density and temperature of the Voronoi cell. The cooling rate and heating rate of the gas cells are updated simultaneously. Crucially, the local chemical state determines the mean molecular weight μ , which is used to calculate the gas temperature T from the internal energy u via the equation of state:

$$T = \frac{(\gamma - 1)u\mu m_p}{k_B}, \quad (15)$$

where m_p is the proton mass, γ is the adiabatic index, u is the specific internal energy, and k_B is the Boltzmann constant. This feedback loop ensures that the collapse of the gas is correctly regulated. For the simulations presented here, we set $\gamma = 5/3$ since the gas is dominated by monatomic particles (H and He in neutral regions or electrons and ions in ionized regions) at all of the densities we resolve.

4.3 Initial conditions

Below we detail how our initial conditions were generated and the motivations for our methods.

4.3.1 Halo generation

In order to isolate the impact of FDM small-scale dynamics from the differences in large-scale structure formation, we initialize our haloes using a non-cosmological, DMO, solitonic core merger

setup, similar to that of P. Mocz et al. (2017). This allows us to generate CDM and FDM haloes with identical global properties, specifically total halo mass, total energy, and angular momentum. In full cosmological volumes, the FDM suppression of small-scale power naturally leads to different halo histories, and therefore makes comparing the gas processes and collapse between haloes of different DM models difficult to disentangle from the effects of different formation histories. In contrast, with our controlled setup we can isolate any effect that the bulk halo properties might have when comparing our FDM, ‘frozen’ FDM, and CDM cases.

For each halo generation (both FDM and CDM), we place a set of solitonic cores (as defined by equation 7) within our box and allow them to merge. We place solitons randomly across the box, excluding the volumes within 8 kpc of the box edges. The mass of each of the solitons is also randomly selected from a uniform distribution over $U(M_h/12, M_h/6)$ such that we place between 6 and 12 solitons in the box. The mass of the final soliton is fixed such that the total mass in the box is M_h . The radius of each soliton is then defined in terms of its mass via the core radius–mass relation of H.-Y. Schive et al. (2014a), P. Mocz et al. (2017):

$$M_c(r_c) = 3.59 \times 10^7 \left(\frac{2.5 \times 10^{-22}}{m_a} \right)^2 \left(\frac{r_c}{\text{kpc}} \right) M_\odot. \quad (16)$$

We illustrate the process of halo generation in Fig. 2. It shows a series of snapshots starting with our initial solitonic cores, and various stages of their merger history, with the final halo shown in the rightmost snapshot. It should be noted that for our simulation boxes, due to the periodic boundary conditions and small overall box size, the virial ratio $2K/|W|$ is greater than one despite the system appearing to reach a steady state. This is likely due to the underestimation of $|W|$ as has been noted in small periodic boxes before (C. Power & A. Knebe 2006; I. Álvarez-Ríos, F. S. Guzmán & P. R. Shapiro 2023; K. Blum et al. 2025).

We evolved these DMO systems for between 1–5 Gyr depending on the halo mass. The initial assembly phase was monitored and stopped when the DM potential has reached a stable, virial-

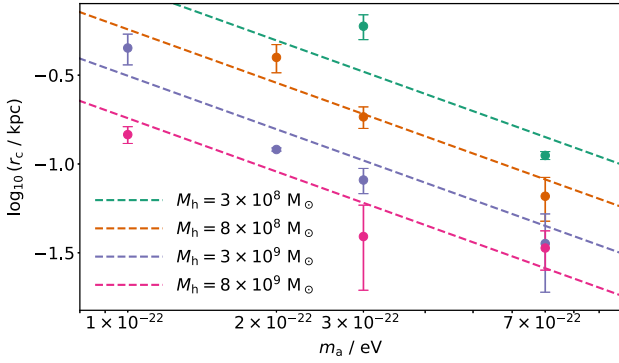


Figure 3. The core radii of our simulated haloes (dots) compared to the core radius–halo mass relation of H.-Y. Schive et al. (2014b, dashed lines). Error bars indicate the 1σ variation in our core radii over time for each halo. While our core radii roughly agree, there is significant variation consistent with the findings of more recent works with large samples of FDM haloes (e.g. B. Schwabe et al. 2016; M. Nori & M. Baldi 2021; H. Y. J. Chan et al. 2022).

ized equilibrium. This ‘settling’ phase is monitored via the energy evolution (Fig. 2); we ensure that the classical kinetic (K_v) energy has plateaued, indicating that the initial ‘ringing’ of the merger has subsided. In FDM haloes, persistent exchange between K_p and W is expected post-virialization due to the large wave effects and the Heisenberg uncertainty principle acting on the wavefunction, see P. Mocz et al. (2017). The lower panel of Fig. 2 shows how the energy components, calculated as in equation (5), change during this merger process, the initial ‘ringing’ of the merger at ~ 235 Myr and subsequent stabilization of K_v . It was only once the initial halo had been assembled that we then stopped the SP solver in order to create our ‘frozen’ FDM haloes.

We also looked at the properties of the solitons generated in our haloes in order to verify consistency with the scaling relations for the soliton core radius found in works like H.-Y. Schive et al. (2014b). In Fig. 3, we show the core radii for our simulated haloes compared to the relation of H.-Y. Schive et al. (2014b, dashed lines). For each of our haloes, we see rough agreement with this derived relation, but with significant variation from halo to halo as was also noted by B. Schwabe, J. C. Niemeyer & J. F. Engels (2016), M. Nori & M. Baldi (2021), and H. Y. J. Chan et al. (2022).

It should be noted that we attempted to run the three combinations of low m_a and low M_h marked with ‘-’ in Table 1 but due to the low mass of these haloes, the periodic boundary conditions, and their low m_a , these setups resulted in diffuse, spatially extended DM distributions that did not form distinct haloes. We also did not run $m_a = 2 \times 10^{-22}$ eV for our highest halo mass for reasons of computational efficiency as we did not expect to see any interesting insights that would not be clear from the $m_a = 1 \times 10^{-22}$ or 3×10^{-22} eV cases.

4.3.2 Initial state of the gas

Once the DM halo is virialized, gas is added as a uniform background, with an initial total of 256^3 gas cells and a total mass consistent with the cosmic baryon fraction. The gas is initially at rest. While a large-scale relative velocity between DM and gas (streaming velocity) of (on average) ~ 30 km s $^{-1}$ is expected at recombination which decays by a factor of $(1+z)$ leading to a relative velocity of $v_{bc} \approx 0.6$ km s $^{-1}$ at $z = 20$ (D. Tseliakhovich

Table 2. Initial chemical species abundances, defined relative to the number of H nuclei. Since we only set the overall deuterium (D) abundance first, then the D^+ and HD abundances, the abundance of all deuterium species together are shown together in the top row.

Species	Initial abundance n_{species}/n_H
D + D^+ + HD	2.6×10^5
H^+	1.0×10^4
H_2	2.0×10^6
H^-	0.0
H_2^+	0.0
D^+	2.6×10^9
HD	0.0
He	7.9×10^2
He^+	0.0
He^{++}	0.0
Free electrons (e^-)	1.0×10^4

& C. Hirata 2010), we omit this effect in this work, as well as any intrinsic turbulence/kinetic energy in the initial gas state, to purely isolate the local impact of FDM dynamics. An initial comparison of the velocity of the soliton’s random walk at the centre of the halo, which is a few tens of km s $^{-1}$, further indicates that any physically meaningful v_{bc} would be negligible in comparison to the effects of the soliton’s dynamics. We hope to address these considerations in future work. The initial gas temperature is set to 10^4 K, representing the typical virial temperature of atomic cooling haloes at $z \sim 20$ (A. T. P. Schauer et al. 2021; M. Kulkarni, E. Visbal & G. L. Bryan 2021; T. Dome et al. 2024). The chemical network is initialized with primordial abundances (Table 2), with a small seed fraction of H_2 (2×10^{-6}).

5 RESULTS

We begin by examining the fundamental properties of the FDM haloes we have generated, which serve as the foundation for all subsequent gravitational effects on the baryons. Unlike the standard CDM paradigm, where haloes are described by a smooth, cusped potential, FDM haloes are characterized by a central solitonic core and persistent, ubiquitous density fluctuations as described in Section 3.2. These fluctuations are a direct consequence of the wave-like nature of the FDM particles, leading to a gravitational potential that is not static but rather constantly in motion. Here, we explore the impact of these fluctuations on the gravitational potential, the collapse of the primordial gas and the subsequent star formation.

5.1 Gravitational potential and density fluctuations at the halo centre

It is well established that the shape of the potential wells of CDM and FDM haloes differ, especially below the core radius r_c . As described in Section 3.1, while CDM haloes are characterized by ‘cuspy’ density profiles at small radii, FDM haloes form a solitonic core at their centre (W. Hu et al. 2000; P. Mocz et al. 2017; S. May & V. Springel 2021; M. Nori, S. Bhatia & A. V. Macciò 2024).

Fig. 4 shows the spherically averaged gravitational potential profile and the RMS fluctuations of the potential for the $M_h = 8 \times 10^8 M_\odot$ haloes, comparing FDM cases with varying m_a to the CDM case. We choose this halo mass as the contract in behaviour between the CDM and lightest m_a FDM cases was most pronounced (See Section 5.5 for further details on the halo mass

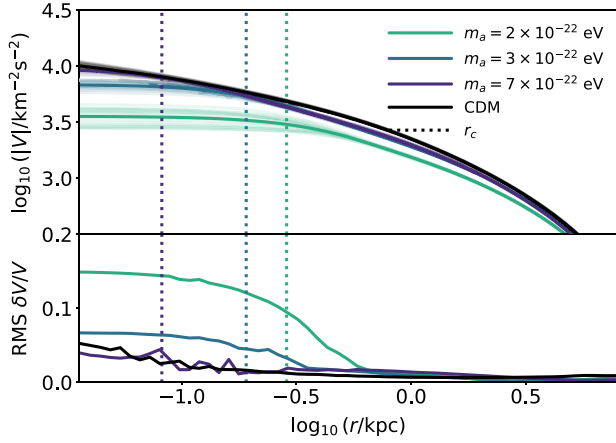


Figure 4. Top: the time-averaged gravitational potential profiles of FDM haloes for different axion masses ($m_a = 2 \times 10^{-22}$, 3×10^{-22} , and 7×10^{-22} eV) and the corresponding CDM halo, all with a mass of $M_h = 8 \times 10^8 M_\odot$. Individual snapshots are shown by the translucent lines. The vertical dotted lines indicate the core radius (r_c) for each FDM halo. The potential flattens inside the core radius for FDM haloes, a signature of quantum pressure. Bottom: the RMS fluctuation of the potential ($\delta V/V$) for the same haloes. This quantifies the fluctuations of the solitonic core at the centre of each FDM halo. The fluctuations are most pronounced and extend to larger radii for smaller m_a , and they decrease significantly outside the solitonic core radius.

dependence of our results). The upper panel displays the time-averaged potential profiles in bold, with the potential profiles for each snapshot shown as translucent lines. For each FDM halo, a core is evident at small radii, where the potential flattens due to the quantum pressure. As the axion mass m_a increases, the FDM potential profile transitions to become more like the CDM profile, which has a sharper cusp. The potential of the $m_a = 7 \times 10^{-22}$ eV halo is very similar to the CDM case, suggesting a reduced impact on gravitational dynamics for higher axion masses.

The lower panel of Fig. 4 shows the dimensionless RMS fluctuation of the potential ($\delta V/V$) as a function of radius. These fluctuations occur on a time-scale of $\sim \tau_{\text{osc}}$ as outlined in Section 3.2 and are a key signature of FDM haloes, resulting from the persistent interference of the DM wavefunction. The plot shows a clear dependence of these fluctuations on the axion mass. For the lightest axion mass $m_a = 2 \times 10^{-22}$ eV that we ran for our $8 \times 10^8 M_\odot$ halo, the RMS fluctuation is notably high at the centre of the halo, reaching a value of ~ 0.15 and extending out to a radius corresponding to the soliton core radius (r_c). As m_a increases, the magnitude of these fluctuations decreases significantly, with the $m_a = 7 \times 10^{-22}$ eV case showing very low variation, comparable to the CDM case. It is worth noting that there are persistent granular fluctuations throughout the halo, but these have a minimal impact on the spherically averaged gravitational potential at radii larger than the soliton radius because they are at much lower density than the soliton itself, and as a result the soliton's properties dominate the shape of the potential well.

The persistent and significant fluctuations in the central gravitational potential, particularly for low m_a , are the primary mechanism through which FDM dynamics affect the baryonic gas. These fluctuations act as a source of energy and turbulence for the gas, hindering its efficient collapse and leading to the effects we discuss in the following subsections.

5.2 Distribution of cold dense gas and FDM dynamics as a driver of turbulence

We now examine how the unique central dynamics of FDM haloes physically restructure the gas distribution. Fig. 5 presents a time-series of gas density snapshots for a halo of $M_h = 8 \times 10^8 M_\odot$, comparing the ‘dynamic’ FDM with $m_a = 2 \times 10^{-22}$ eV to the corresponding ‘frozen’ FDM case and the control CDM scenario. The snapshots are shown starting from the moment of the first sink formation t_* followed by snapshots at $\Delta t \approx 50$, ≈ 100 , and ≈ 190 Myr since t_* . We recall that sink particles, shown in grey in Fig. 5, represent gas that has entered irreversible runaway collapse (see Section 4.2.4), and thus provide a proxy for star formation efficiency.

While the CDM case (right column) exhibits a rapid, centrally concentrated collapse into a single dense region, both FDM cases show a more extended gas morphology. However, a critical distinction emerges between the two FDM models: in the ‘dynamic’ case (left column), the gas distribution is significantly more extended and fragmented than in the ‘frozen’ case (middle column). Although the ‘frozen’ model possesses the same flattened gravitational potential (the soliton core) as the full FDM scenario, the gas settles into a relatively symmetric, albeit wider than in the CDM case, configuration. In contrast, the ‘dynamic’ FDM potential exerts stochastic gravitational torques that continuously stir the gas. This dynamic stirring drives turbulence that counteracts the inward gravitational pull, maintaining a diffuse cloud that becomes increasingly spatially extended over time. Because our simulations employ a strict Jeans refinement criterion that resolves the local Jeans length by at least 16 cells down to the sink formation threshold, the formation of multiple independent sink particles is a genuine physical signature of gas fragmentation rather than numerical instability.

This physical disruption is further quantified by the evolution of the sink particle half-mass radius, represented by the red circles in Fig. 5. This half-mass radius is the radius of a sphere around the centre of mass of sinks that contains half of the total mass in sinks. In the CDM halo, sinks form a compact central cluster with a stable, small half-mass radius, indicating efficient and localized gas collapse. Conversely, in ‘dynamic’ FDM, the sinks are found far from the halo centre; they are distributed across a wider volume, and their half-mass radii often increase as the simulation progresses. This dramatic increase in the half-mass radius is a direct signature of the FDM soliton’s dynamics.

The fact that the half-mass radius in the ‘frozen’ case remains significantly smaller than in the ‘dynamic’ case (though larger than CDM) confirms that while the core geometry sets a floor for the spatial extent of star formation, the fluctuation dynamics are the primary driver of the observed distribution. This implies that the spatial clustering of the first stars in FDM cosmologies may be much lower than predicted by static halo models, potentially impacting the local feedback, binary formation, and subsequent galaxy assembly.

The impact of these dynamics is further quantified in Fig. 6, which displays the spherically averaged radial density profiles for DM, gas, and sinks in the $8 \times 10^8 M_\odot$ halo. To account for the highly dynamic and time-variable nature of the gas and sink distribution within our simulations, the gas and sink particle profiles shown are averaged over five late-time snapshots centred around $t \approx 250$ Myr from the start of the simulations. In agreement with literature, our DM profiles (top panel), which have been normalized by the core radius r_c , clearly distinguish the FDM models

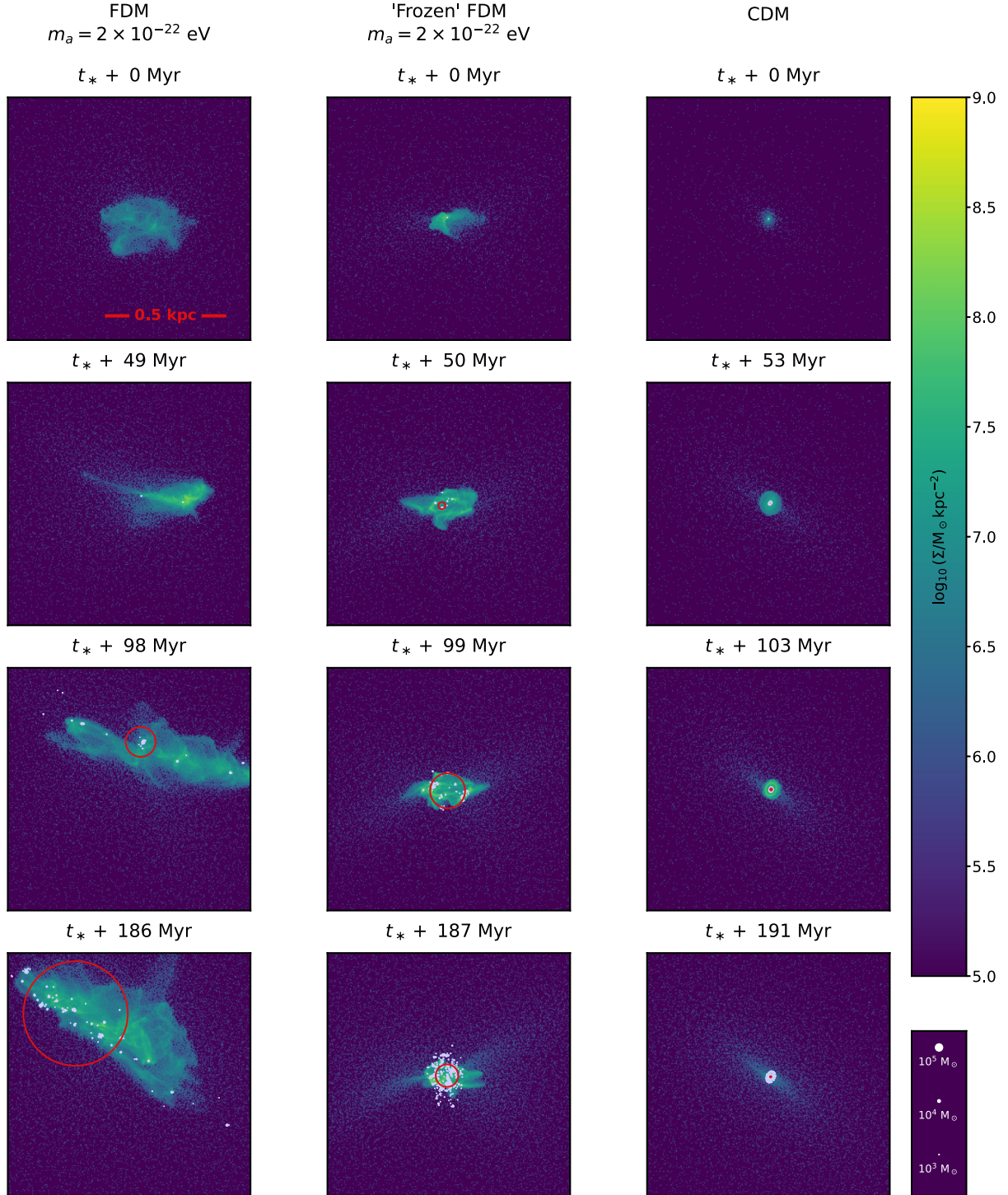


Figure 5. Gas density projections in the central region of each halo for a selection of our simulations with a halo mass of $8 \times 10^8 M_\odot$: FDM with $m_a = 2 \times 10^{-22}$ eV (left), the ‘frozen’ FDM with the same $m_a = 2 \times 10^{-22}$ eV (centre), and a CDM case (right). The first row shows the gas state at the formation of the first sink. As the simulation proceeds, in the FDM cases, the gas is distributed more widely across the central region of the halo as it is affected by the dynamics of the central soliton. In the CDM case, gas remains tightly bound in the central region. The dots show the location of the sink particles, with the radius scaled by the sink mass $m_{\text{sink}}^{1/3}$. The red circle shows the half-mass radius of the sink particles. There is no half mass radius shown for the ‘dynamic’ FDM case for $t_* + 49$ Myr as there was still only one sink formed at this time.

from the CDM cuspy central density by the presence of a central solitonic core.

The gas density profiles (middle panel) reveal a complex relationship between the DM potential and the efficiency of gas

collapse. In the CDM case, the central gas density is lower than in the FDM models up to very small radii $r \lesssim 10^{-1.5} r_c$. It is only below this small radius where the gas density increases above the level seen in the FDM haloes. This is due to the highly efficient

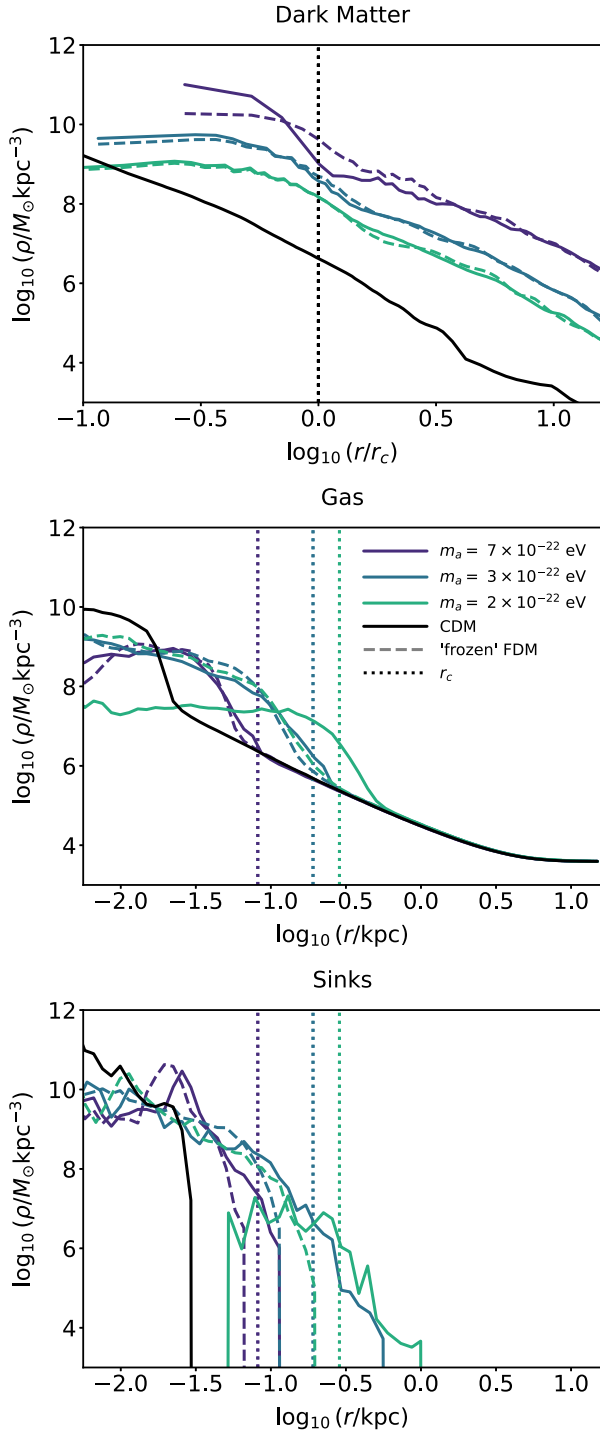


Figure 6. Radial density profiles for the $8 \times 10^8 M_\odot$ halo at a late-time snapshot with $t \approx 250$ Myr from the start of the simulations. The dashed lines show the ‘frozen’ FDM cases, while the dotted lines show the core radius r_c . Top: DM density profiles normalized by the core radius r_c . For completeness, we show CDM normalized with $r_c = 1$ kpc. The FDM cases exhibit a flattened solitonic core, while the CDM shows a characteristic cusp. Middle: gas density profiles averaged over five late-time snapshots. Bottom: radial distribution of sink particles averaged over five late time snapshots. In the FDM simulations, particularly at lower axion masses, the sinks are spatially dispersed and displaced from the halo centre. The vertical dotted lines indicate the core radius r_c for each respective FDM model.

accretion of gas on to the few sink particles that form at the centre of the halo, which depletes the central gas reservoir. In the FDM cases, however, the gas more closely tracks the shape of the solitonic core and remains at high densities out to larger radii as m_a is decreased. This suggests that while the gas is able to accumulate within the core, in the FDM case the gas is prevented from completing the transition into runaway collapse into sinks as efficiently as in CDM. This is further discussed in Section 5.5.

A comparison between the FDM models shows that the central gas density is suppressed for the lowest axion masses, with the $m_a = 2 \times 10^{-22}$ eV case showing a lower central peak than either the $m_a = 7 \times 10^{-22}$ and 3×10^{-22} eV or any of the ‘frozen’ models. This confirms that the injection of energy and turbulence by the halo dynamics is more pronounced at lower axion masses, and actively inhibits the gas from reaching the densities required for sink formation for the lowest axion masses, even when a substantial gas reservoir is present.

Finally, the radial distribution of sink particles (bottom panel) illustrates the spatial displacement of the sinks from the halo centre. The sinks in the ‘dynamic’ FDM models are scattered across a wider volume when compared to their ‘frozen’ counterparts, particularly for $m_a = 2 \times 10^{-22}$ eV. This larger scale distribution is a hallmark of the FDM dynamics; as the gravitational potential minimum shifts stochastically, the sites of gas collapse are never fixed, leading to a fragmented and non-centralized distribution of the sinks within our simulations.

5.3 Angular momentum and rotational support

We might expect that the dynamics at the centre of the FDM haloes can induce torques on the gas if the collapse time-scale is longer than that of the dynamical time-scale of the FDM fluctuations, and hence lead to a degree of enhanced rotational support against collapse. This can be quantified by examining the specific angular momentum (j) of the baryonic matter (gas plus sinks) as a function of radius in each of our haloes. We can compare this to the specific angular momentum required to support the gas against collapse at a given radius. For any given mass element, its specific angular momentum around a given axis is given by $\mathbf{j} = \mathbf{r} \times \mathbf{v}$. We can then link this to a centripetal force at a given radius r via

$$F_{\text{cent}} = \frac{v_\phi^2}{r} = \frac{j(r)^2}{r^3}, \quad (17)$$

where v_ϕ is the velocity perpendicular to the radial vector, and $j(r) = |\mathbf{j}(r)| = rv_\phi$ is the magnitude of the total specific angular momentum of a spherical shell at radius r . We can then link this force to the gravitational force at a given radius to find an estimate for the angular momentum required to support the gas as a function of radius via $F_{\text{cent}} = -F_{\text{grav}} = \nabla V$, thus arriving at

$$\frac{j_{\text{crit}}(r)^2}{r^3} \approx \frac{\partial V}{\partial r} \Rightarrow j_{\text{crit}}(r) \approx \sqrt{\frac{\partial V}{\partial r}} r^3. \quad (18)$$

Fig. 7 shows the specific angular momentum j calculated for each spherical shell for the combined gas and sinks as a function of r from the halo centre for both CDM and various FDM cases, alongside the critical angular momentum, j_{crit} , required for rotational stabilization against collapse (shown with dot-dashed lines in Fig. 7). The key feature of these data is the significant amount of rotation induced within the gas in our ‘dynamic’ FDM haloes, particularly for radii below the soliton radius, r_c , and for lighter axion masses, m_a .

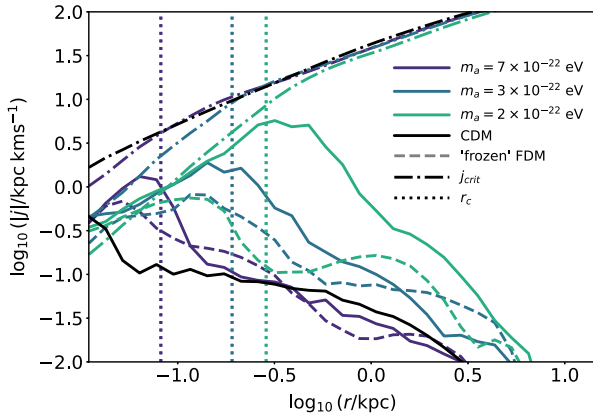


Figure 7. Time-averaged specific angular momentum (j) of gas and sinks as a function of radius (r) from the halo centre. The solid black line represents the CDM case, solid coloured lines represent FDM cases with different axion masses (m_a , see the legend), while dashed lines correspond to the ‘frozen’ FDM cases. The dash-dotted lines indicate the critical angular momentum (j_{crit}) required for rotational stabilization against gravitational collapse. Vertical dotted lines denote the empirical value for the soliton radius (r_c) for each FDM halo mass (H.-Y. Schive et al. 2014b). Note the significant induced rotation in FDM cases (especially for lower m_a) below r_c , where j approaches or exceeds j_{crit} , in contrast to the minimal induced rotation in the CDM and the ‘frozen’ FDM cases. All data are taken from the $8 \times 10^8 M_\odot$ halo mass case at a late evolutionary time, with the snapshots taken from more than 100 Myr after the formation of the first sink.

For the ‘dynamic’ FDM cases, especially $m_a = 2 \times 10^{-22}$ and 3×10^{-22} eV, the specific angular momentum j of the gas within the central regions ($r \lesssim r_c$) is notably higher compared to the CDM and ‘frozen’ FDM counterparts (as shown in Fig. 7). This trend suggests that the induced rotation in FDM cases is a direct consequence of the dynamic nature of the solitonic core. The persistent random walk and density fluctuations of the FDM soliton exert time-varying gravitational torques on the surrounding gas, imparting angular momentum to it. This effect is more pronounced for lighter axion masses because the de Broglie wavelength is larger, leading to more extended and hence gravitationally influential solitonic cores and a greater effect on the gas out to larger radii. Furthermore, the amplitude of these density fluctuations also scales with λ_{dB} , as the larger coherence length in lighter axions results in more massive, slower moving interference granules that exert stronger and more persistent influence on the solitonic core.

Crucially, in the ‘dynamic’ FDM cases, this specific angular momentum reaches or exceeds the critical threshold (j_{crit}) required to rotationally stabilize the gas against collapse (dash-dotted lines, Fig. 7). While the ‘frozen’ FDM results (Section 5.5) demonstrate that a static cored potential is sufficient to delay the initial timing of gas collapse, the dynamic treatment reveals a secondary, structural effect: the gas is prevented from reaching the same high-density, compact configurations seen in the static cases. Instead, the injection of angular momentum and turbulence leads to a more spatially extended gas distribution. In contrast, the ‘frozen’ FDM and CDM cases show negligible induced rotation, allowing gas to settle more efficiently into the potential minimum. This highlights a key physical distinction: whereas the soliton core geometry sets the timing of star formation, the soliton dynamics dictate the spatial scale of the resulting dense gas

structures. By identifying this rotational stabilization (a mechanism currently absent from analytical models of FDM-driven star formation), we suggest that the suppression of gas efficiency in FDM haloes is likely greater than previously estimated.

5.4 Energy budget

To summarize the mechanisms preventing efficient gas collapse and driving the observed morphology, we examine the energy budget of the gas. Fig. 8 shows the radial profiles of the specific energy components (top row) alongside a detailed decomposition of the specific kinetic energy into its radial (K_{rad}), rotational (K_{rot}), and turbulent (K_{turb}) modes (bottom row), comparing the CDM, ‘frozen’ FDM, and ‘dynamic’ FDM scenarios for the $M_h = 8 \times 10^8 M_\odot$ halo. We decompose the total specific kinetic energy of the gas within each radial shell into three components:

$$K_{\text{tot}} = K_{\text{rad}} + K_{\text{rot}} + K_{\text{turb}}. \quad (19)$$

Here, the specific radial kinetic energy is given by:

$$K_{\text{rad}} = \frac{1}{2} \frac{1}{M_{\text{shell}}} \sum_i m_i (\mathbf{v}_i \cdot \hat{\mathbf{r}}_i)^2, \quad (20)$$

where $\hat{\mathbf{r}}_i$ is a unit radial vector and $M_{\text{shell}} = \sum m_i$ within each spherical shell. To determine the rotational component, we use the net specific angular momentum vector of each individual shell to define an axis of rotation, and then calculate the specific rotational kinetic energy K_{rot} similarly to above, but taking the vector cross product of the velocities with a unit vector aligned with the axis of rotation. Any residual kinetic energy not captured by these ordered components is attributed to gas turbulence, defined simply as $K_{\text{turb}} = K_{\text{tot}} - K_{\text{rad}} - K_{\text{rot}}$.

There are clear differences in the thermal evolution of the gas between the cosmologies. In the CDM case (top-left panel), the specific thermal energy (U) remains remarkably high down to the smallest resolved radii, balancing the deep gravitational potential well as the gas undergoes concentrated gravitational compression. Conversely, inside both FDM haloes, a major transition occurs near the core radius r_c . Below r_c , the gas undergoes significant cooling, causing the thermal energy to drop by an order of magnitude. Without sufficient thermal pressure to support the gas, the prevention of collapse is likely mechanical rather than thermal.

The underlying nature of this mechanical support is clearly illustrated by the kinetic energy breakdown (bottom row). In the CDM halo, the kinetic energy budget is overwhelmingly dominated by the radial component ($K_{\text{rad}}/K_{\text{tot}} \gtrsim 0.85$) across all radii, reinforcing the interpretation of coherent, uninhibited inflow directly into the central potential minimum. In contrast, both the ‘frozen’ and ‘dynamic’ FDM cases exhibit a dramatic redistribution of kinetic energy into non-radial modes beneath the solitonic core radius. For $r < r_c$, non-radial support becomes dominant, forcing the fractional radial kinetic energy down to less than 30 per cent of the total kinetic energy budget.

Crucially, the mechanism generating this non-radial kinetic energy differs between the two FDM treatments, highlighting the role of active soliton fluctuations. In the ‘frozen’ FDM case (bottom-middle panel), the suppressed radial kinetic energy is converted largely into turbulent modes (K_{turb}), while the induced rotational component (K_{rot}) remains low. This transition is bounded by the core radius, with the gas outside r_c remaining predominantly radial.

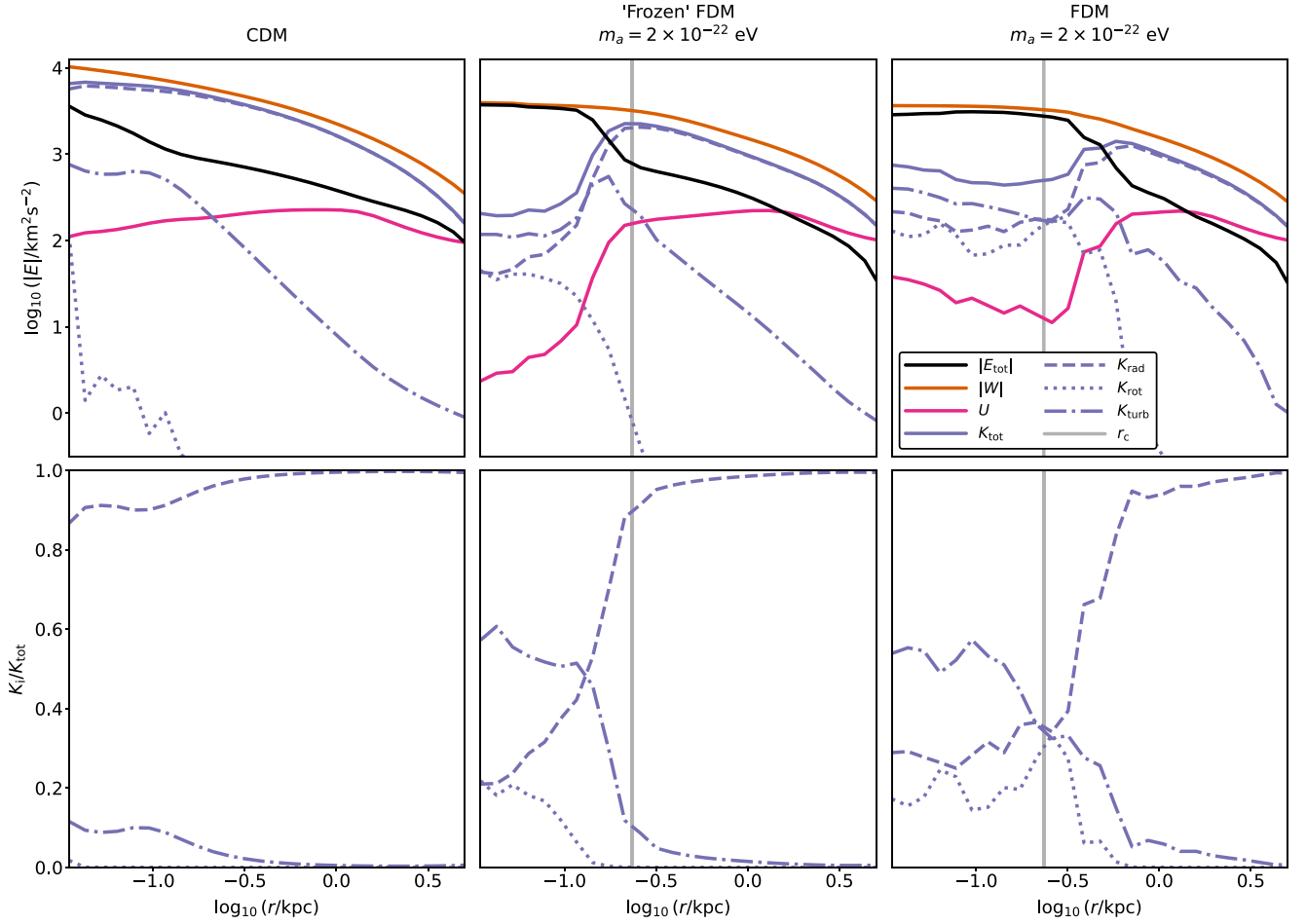


Figure 8. The energy budget of the gas for our $M_h = 8 \times 10^8 M_\odot$ halo for CDM (left) and FDM with $m_a = 10^{-22}$ eV for our ‘frozen’ (middle) and ‘dynamic’ (right) setups. The top row shows the absolute value of each energy component: thermal (U), kinetic (K), potential (W), and the total (E). The bottom row shows the fraction of kinetic energy in radial (K_{rad} , dashed), turbulent (K_{turb} , dot-dashed) and rotational (K_{rot} , dotted) modes.

In contrast, the ‘dynamic’ FDM potential (bottom-right panel) acts as an efficient engine for both turbulence and large-scale rotation. It induces a boost in rotational kinetic energy (K_{rot}) out to radii at and beyond the r_c that is absent in the static cored potential. Likewise, the ‘dynamic’ FDM is able to actively pump turbulence beyond r_c , unlike the ‘frozen’ case. This extended injection of mechanical energy explains why the ‘dynamic’ gas distribution remains significantly more diffuse and spatially extended than its ‘frozen’ counterpart, establishing that the time-dependent fluctuations of the FDM haloes are central to reshaping the environment of the first stars.

5.5 Accumulation of gas in sinks: core versus dynamics

Having looked qualitatively at how the gas is affected by the FDM dynamics, we now quantify the effects on the total accumulation of cold, dense gas. As detailed in Section 4.2.4, we employ sink particles as a numerical proxy for gas that has entered the runaway collapse phase ($n > 10^2 \text{ cm}^{-3}$) representing pre-stellar collapse. This allows us to track the prospects of each halo to form stars and to make comparisons between haloes. Fig. 9 illustrates the cumulative total mass in sinks $M_{\text{sink}} = \sum_i m_{\text{sink},i}$ (where $m_{\text{sink},i}$ is the mass of an individual sink in a given simulation) over time across our halo suite, providing a direct measure of the relative star formation efficiency in FDM versus CDM environments.

5.5.1 High-mass regime: dominance of halo geometry

In our most massive haloes ($M_h = 8 \times 10^9$ and $3 \times 10^9 M_\odot$), the primary driver of suppression of gas in sinks is the cored geometry of the FDM potential. As shown in Fig. 9, the mass accumulated in the ‘frozen’ FDM cases (dashed lines) closely tracks the fully ‘dynamic’ FDM runs (solid lines).

For our heaviest halo $8 \times 10^9 M_\odot$, we observe a modest 11 per cent reduction in total sink mass for the lightest simulated FDM case, $m_a = 10^{-22}$ eV, compared to CDM. In this regime, the gravitational potential is sufficiently deep that the stochastic fluctuations of the soliton act as a subdominant perturbation. Instead, the primary suppression comes from the flattened central potential gradient (the solitonic core), which reduces the gravitational acceleration in the central region, lowering the efficiency of gas inflow. After a brief initial delay in the onset of star formation, the sink accumulation rate reaches an approximately steady state relative to CDM, indicated by the flattening of the ratio in the lower subpanels.

5.5.2 Low-mass regime: dominance of FDM dynamics

The impact of FDM physics becomes dramatically more pronounced in the lower mass haloes ($M_h = 8 \times 10^8$ and $3 \times$

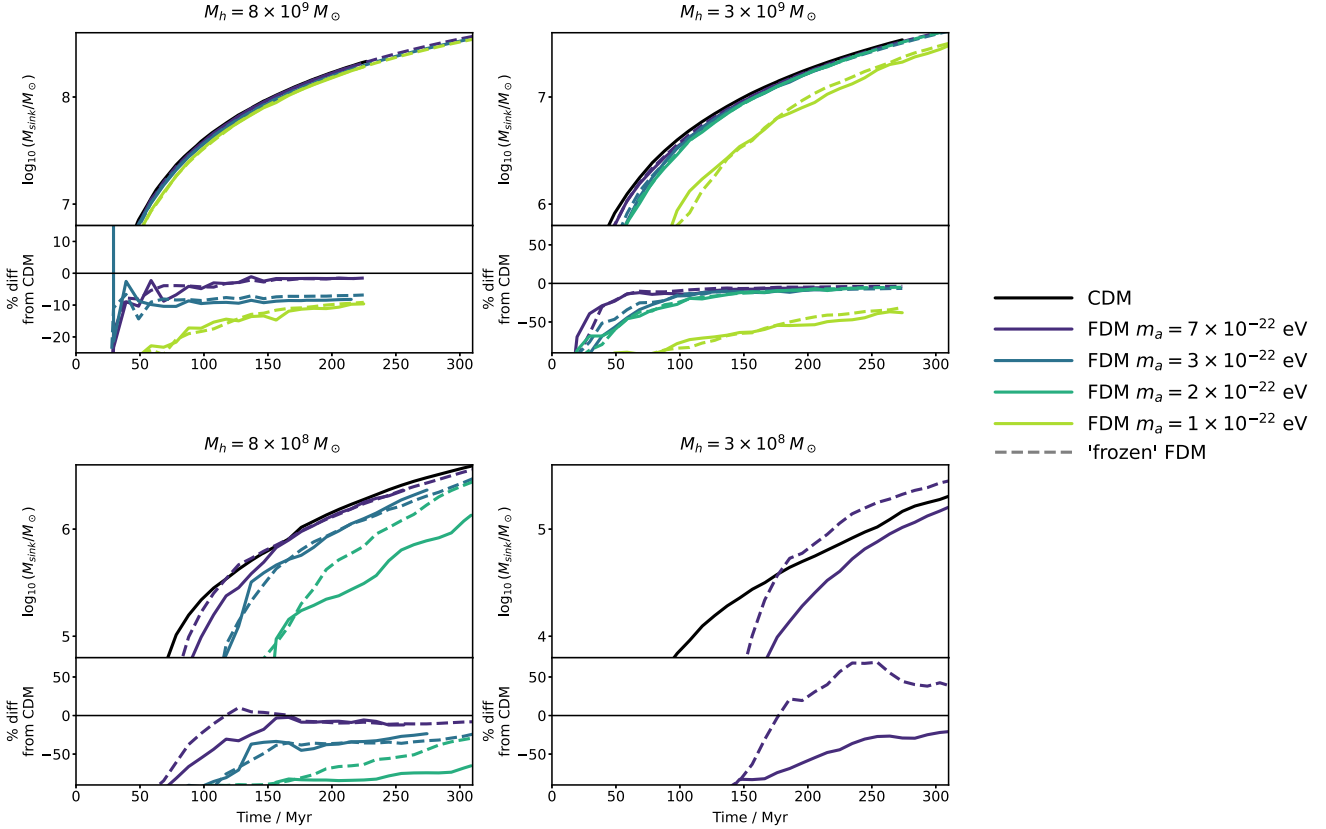


Figure 9. Cumulative total mass in sinks $M_{\text{sink}} = \sum_i m_{\text{sink},i}$ (where $m_{\text{sink},i}$ is the mass of an individual sink in a given simulation) over time across the four simulated halo masses (arranged in decreasing mass from top-left to bottom-right). Each of the four panels is divided into an upper plot showing the absolute total sink mass, and a corresponding lower plot displaying the percentage difference in mass accumulation relative to the CDM baseline. Dashed lines represent the ‘frozen’ FDM cases.

$10^8 M_{\odot}$), where the DM dynamics emerge as the dominant suppression mechanism.

For the $M_h = 8 \times 10^8 M_{\odot}$ case, we find up to an 84 per cent reduction in sink mass for the ‘dynamic’ run compared to CDM obtained for the lowest axion mass we were able to simulate in this halo, $m_a = 2 \times 10^{-22}$ eV. Crucially, the ‘frozen’ FDM run for this same halo shows a significantly higher accumulation rate than its ‘dynamic’ counterpart. This gap confirms that core geometry alone cannot account for the suppression, and the time-varying dynamics of the fully simulated FDM halo actively prevent gas from settling into the sink-forming phase.

In our lowest mass haloes ($M_h = 3 \times 10^8 M_{\odot}$), star formation is almost entirely inhibited for lighter axion masses. Only the $m_a = 7 \times 10^{-22}$ eV run formed sinks within the 300 Myr window, showing a substantial (23 per cent) but inconsistent reduction compared to CDM. For the other axion mass run that we performed for this halo mass ($m_a = 3 \times 10^{-22}$ eV), no sinks were formed within the first 300 Myr (the first sinks formed more than 500 Myr after the start of the simulation). This near-total suppression occurs because the de Broglie wavelength and hence soliton radius is comparable to the virial radius of the halo. In this state, the soliton fluctuations disrupt the entire central gas region, preventing the formation of a bound, cooling core. The percentage suppression in total sink mass for three late times (200, 250, and 300 Myr after the start of the simulation) are summarized for all cases in Table 3.

Interestingly, in our lowest-mass halo ($M_h = 3 \times 10^8 M_{\odot}$), the ‘frozen’ FDM model ($m_a = 7 \times 10^{-22}$ eV) eventually surpasses the CDM case in total sink mass by nearly 61 per cent at $t = 250$ Myr, despite experiencing a significant initial delay (see Table 4). This was the only FDM run in our suite to show such behaviour. We interpret this as the effect of a massive, static soliton core. In this lowest mass regime, the soliton constitutes a dominant fraction of the total halo mass. When this mass is held static (as in the ‘frozen’ case), it creates a wide, deep, and perfectly stable potential well that, once gas cooling is initiated, facilitates a more efficient and monolithic accumulation of gas than the tidally active CDM cusp.

Crucially, this behaviour disappears entirely in the ‘dynamic’ FDM run, where the sink mass is suppressed by 58 per cent at 200 Myr and remains significantly lower than CDM throughout. This stark contrast between the ‘frozen’ and ‘dynamic’ results at $3 \times 10^8 M_{\odot}$ provides evidence that the suppression of star-forming gas in low-mass FDM haloes is not a geometric effect of the cored potential, but is driven by the kinetic energy and angular momentum injected by the stochastic soliton fluctuations.

5.5.3 Calibrating semi-analytical models

The observed flattening of the sink mass ratios at later times for our higher halo masses ($M_h = 8 \times 10^9$ and $3 \times 10^9 M_{\odot}$) provides a vital calibration point for semi-analytical star formation

Table 3. Evolution of percentage reduction in accumulated sink mass relative to CDM. Percentages are calculated at three late-time snapshots ($t = 200, 250,$ and 300 Myr) after gas introduction. For our higher halo masses, the later times are not shown as the simulations did not run to that time due to computational time constraints. A value of -100 per cent indicates total suppression (no sinks formed), while positive values (e.g. the $3 \times 10^8 M_\odot$, $m_a = 7 \times 10^{-22}$ eV ‘frozen’ case) indicate an increase in mass relative to CDM.

Halo mass ($M_\odot$)	Time (Myr)	Mass reduction relative to CDM (per cent)							
		$m_a = 10^{-22}$ eV		$m_a = 2 \times 10^{-22}$ eV		$m_a = 3 \times 10^{-22}$ eV		$m_a = 7 \times 10^{-22}$ eV	
		Dynamic	Frozen	Dynamic	Frozen	Dynamic	Frozen	Dynamic	Frozen
8×10^9	200	−11	−10	−	−	−8	−7	−2	−2
3×10^9	200	−54	−45	−8	−8	−9	−8	−6	−5
	250	−42	−36	−6	−6	−8	−7	−6	−4
8×10^8	200	−	−	−84	−68	−37	−36	−8	−9
	250	−	−	−76	−54	−28	−35	−12	−11
	300	−	−	−70	−32	−	−28	−	−9
3×10^8	200	−	−	−	−	−100	−100	−58	+25
	250	−	−	−	−	−100	−100	−29	+61
	300	−	−	−	−	−100	−100	−23	+41

Note. Entries marked ‘−’ indicate that the simulation did not reach that time, or was not run entirely (as outlined in Table 1).

Table 4. Time delay of first sink formation (t_{delay}) in FDM haloes relative to CDM. In CDM, sink formation starts at 21, 12, 24, and 60 Myr (from the most massive halo to the lightest, in the order presented in the first row of the table) from the moment the gas is introduced into the virialized halo.

Halo Mass ($M_\odot$)	Time for formation of first sink t_{delay} (Myr)							
	$m_a = 1 \times 10^{-22}$ eV		$m_a = 2 \times 10^{-22}$ eV		$m_a = 3 \times 10^{-22}$ eV		$m_a = 7 \times 10^{-22}$ eV	
	Dynamic	Frozen	Dynamic	Frozen	Dynamic	Frozen	Dynamic	Frozen
8×10^9	0	1	−	−	−2	0	1	1
3×10^9	18	50	6	10	13	6	6	7
8×10^8	−	−	89	63	55	59	37	31
3×10^8	−	−	−	−	631	473	79	75

Note. Cells marked ‘−’ indicate parameter combinations that were not simulated.

models. Our results suggest that while star formation eventually occurs in FDM haloes, as a result of the mechanisms described in Sections 5.3–5.5 the efficiency is severely reduced for m_a at the bottom end of the range explored in this work for the low halo mass end of the FDM halo mass function of any FDM cosmology. Since the sink accumulation rate tracks the availability of cold, collapsing gas, we expect this suppression to translate into a proportionally reduced star formation rate. This implies that the transition from a dark, non-star-forming halo to a star-forming one is markedly different in FDM compared to CDM. We expect this change in star formation efficiency and distribution to significantly affect the predicted high-redshift UVLFs and the timing and duration of reionization (R. Barkana & A. Loeb 2001; P. Dayal & A. Ferrara 2018; J. Mirocha & S. R. Furlanetto 2019). However, the exact form of this efficiency scaling is difficult to determine from the sparse sampling of m_a and M_h that we have in this work.

5.6 Delay in first sink formation

Our results indicate that the formation of the first sink is often delayed in FDM compared to CDM, even for haloes of the same mass, with the magnitude of the delay scaling inversely with both halo and axion masses. These delays in the formation of the first sink are shown in Table 4. In the high-mass regime ($M_h = 8 \times 10^9 M_\odot$), the delay is negligible ($|t_{\text{delay}}| \lesssim 2$ Myr across all models), as the deep gravitational potential quickly overcomes the pressure support of the gas even in the FDM haloes.

However, in the low-mass regime ($M_h \leq 8 \times 10^8 M_\odot$), the delay becomes much more significant across all the FDM cases when compared to the counterpart CDM halo. For these halo masses we see at least a 30 Myr gap between the first CDM sink

formation and all of the FDM haloes. This delay is important to include when calibrating models of star formation in FDM cosmologies to the corresponding CDM cases, especially at high redshifts where a few tens of Myr is a comparable time-scale to the lifetimes of the first population of stars. A delay of ~ 100 Myr in the emergence of the first stellar population would have a detectable impact on the timing of reionization (R. Barkana & A. Loeb 2001; B. E. Robertson et al. 2015) and the structure of the 21-cm signal (e.g. A. Fialkov et al. 2014).

By comparing the ‘frozen’ and ‘dynamic’ columns in Table 4, we can isolate the physical drivers of this delay. For $M_h \gtrsim 3 \times 10^9 M_\odot$, the delay times are similar for the ‘frozen’ and ‘dynamic’ cases, suggesting that the flattened potential gradient of the cored FDM halo profiles is the primary cause of the delay. In contrast, for our lighter haloes ($M_h \leq 8 \times 10^8 M_\odot$) and the lightest axion masses, the ‘dynamic’ FDM run experiences a greater delay in first sink formation. For the lightest halo ($M_h = 3 \times 10^8 M_\odot$) and with $m_a = 3 \times 10^{-22}$ eV, the ‘dynamic’ FDM run is delayed by an additional ~ 160 Myr relative to the ‘frozen’ case. This confirms that at small scales $r \lesssim r_c$, the soliton fluctuations provide a secondary, dynamical barrier to collapse, likely through the injection of turbulent energy as hypothesized in Section 5.2.

To our knowledge, this is the first work to explicitly separate the effect of the geometric delay (core-driven) from the dynamical delay (fluctuation-driven) in a controlled simulation suite. While previous studies have noted a general delay in FDM star formation due to the power spectrum cut-off (e.g. S. Hirano et al. 2018; P. Mocz et al. 2019; D. Jones et al. 2021), our results confirm the importance of the dynamics at the halo centre, which, for the low M_h and m_a end of our parameter space, significantly extends the non-star-forming era of the smallest haloes beyond what a static cored-potential model would predict.

5.7 Evolution of sink particles

Tracing the individual motion and growth of star-forming regions (using sink particles as a proxy) provides a direct probe of the small-scale dynamics within the DM potential. It is important to note that at our resolution, these sink particles represent dense, gravitationally bound gas clouds with the potential to form clusters of stars, rather than individual stellar objects. Consequently, the sink mass and the resulting mass function should be interpreted as indicators of the total reservoir of star-forming gas in a localized region rather than a direct prediction of the stellar initial mass function (IMF). In this section, we examine the structural evolution of the sink population, focusing on the $M_h = 3 \times 10^9 M_\odot$ halo where the transition from core-dominated to dynamics-influenced behaviour is most notable. Fig. 10 displays the number of sink particles in the mass-radius (M_{sink} versus r) plane across three characteristic times: early (close to the first sink formation, 100 Myr), intermediate (150 Myr), and late (250 Myr).

5.7.1 CDM: centrally concentrated accretion

In the CDM case (top row, Fig. 10), the sink population follows the established picture of hierarchical collapse. Sinks typically form very close to the halo centre at $r < 0.04$ kpc and remain there as the population of sinks have no means of gaining energy.

The vertical tracks in the $M_{\text{sink}}-r$ plane indicate that CDM sinks are characterized by sustained accretion. As they settle into the stable potential minimum, they are continuously fed by accreting gas. This results in a ‘top-heavy’ sink mass function (Fig. 11), where the total number of sinks remains relatively stable increasing from $N = 133$ at 100 Myr to $N = 144$ at 250 Myr. However, the median sink mass grows significantly from $M_{\text{med}} \approx 3 \times 10^4 M_\odot$ at 100 Myr to $2 \times 10^5 M_\odot$ at 250 Myr. This behaviour is consistent with the monolithic collapse characteristic of standard primordial star formation (e.g. V. Bromm 2013).

5.7.2 FDM: dynamical scattering and starvation

The evolution of the sink populations in FDM haloes is fundamentally different. While both the ‘frozen’ and ‘dynamic’ FDM cases show a delay in formation, their subsequent behaviours diverge based on the nature of the potential.

In the ‘dynamic’ FDM case (bottom row, Fig. 10), we observe a significant migration of sinks to larger radii. Unlike the CDM sinks, which remain close to the halo centre, sinks in the ‘dynamic’ FDM haloes often move outwards to larger radii comparable to the core radius. As the soliton core undergoes its stochastic random walk, the site of the potential minimum shifts on time-scales comparable to the gas cooling time. This leads to a population of sinks formed near the soliton radius ($r \approx r_c$) that are subsequently ejected to larger radii by the time-varying gravitational potential.

Crucially, once these sinks are scattered away from the dense central gas reservoir, their accretion ceases. This gas starvation results in a sink mass function that is markedly different from CDM: rather than a few massive sinks, the FDM halo hosts a growing number of low-mass clumps (increasing from $N = 154$ at 100 Myr to $N = 2801$ at 250 Myr). The median mass remains stagnant near the formation threshold ($M_{\text{med}} \approx 4 \times 10^3 M_\odot$), as can be seen by the lack of vertical growth in the bottom row of Fig. 10, and in the much more bottom-heavy sink mass function in Fig. 11.

In the ‘frozen’ FDM case (middle row), we see an intermediate behaviour. While the sinks form at larger radii than in CDM due to the flattened core potential, they lack the extreme outward scattering seen in the ‘dynamic’ case. They tend to settle slowly toward the static gravitational centre, but their accretion is still less efficient than in CDM because the core’s lower density reduces the gas supply rate.

This evidence highlights that small FDM haloes may harbour an altered population of the first stars. By scattering star-forming clumps out of high-density regions, FDM dynamics effectively limit the maximum mass any single star-forming region can attain, potentially shifting the IMF of Pop III stars toward lower masses.

The spatial distribution of these formation events and the subsequent growth history of individual sinks are further illustrated in Appendix A, where we show the sink population in radius-mass space. This visualization confirms that FDM sinks not only form at larger radii on average but see a general migration outwards in the halo.

5.8 Chemical stirring in the central halo region

While the initial collapse of gas into the virialized haloes in our mass range (10^8 – $10^9 M_\odot$) is dominated by atomic hydrogen (Lyman- α) cooling, the subsequent formation of stars is contingent upon the production of molecular hydrogen (H_2). Efficient cooling to the temperatures needed for fragmentation ($T \lesssim 500$ K) requires the presence of H_2 as the primary coolant (T. Abel et al. 2002; V. Bromm & R. B. Larson 2004; S. C. O. Glover & T. Abel 2008; V. Bromm & N. Yoshida 2011). In this section, we examine how FDM dynamics alter the spatial and phase-space distribution of this crucial molecule, which has direct implications for the fragmentation and eventual IMF of the first stars.

Fig. 12 shows the state of the gas for a $M_h = 8 \times 10^8 M_\odot$ halo, comparing the ‘dynamic’ FDM with $m_a = 2 \times 10^{-22}$ eV, the corresponding ‘frozen’ FDM, and CDM cases. We synchronize these snapshots with the final epoch of our density series (Fig. 5), though they represent a slice through the halo centre, rather than a projected density as in Fig. 5, to resolve the local chemical variations.

5.8.1 FDM stirring of H_2 -enriched gas

In the ‘dynamic’ FDM case, H_2 -rich gas is liberated from the central potential. In standard CDM (right column), high H_2 fractions ($n_{H_2}/n_H > 10^{-3}$) are strictly confined to the highest density regions ($n \gtrsim 10 \text{ cm}^{-3}$) at the very centre of the halo. In contrast, the ‘dynamic’ FDM case exhibits plumes of H_2 -rich gas extending far into the lower density outskirts of the halo, as can be seen in the top row of Fig. 12.

We use the term ‘liberated’ to describe gas that was processed in the high-density core, where H_2 formation is catalyzed, and subsequently transported to lower density environments. This transport is driven by the stochastic ‘stirring’ of the solitonic core’s random walk. As the potential minimum shifts, gas that was once at the centre is left in lower density regions but retains its high H_2 fraction. Because the cooling time is finite, this gas remains significantly colder than its surroundings, as evidenced by the middle row of Fig. 12, where the $T < 10^3$ K regions overlap almost perfectly with the redistributed H_2 .

This dynamical mixing is quantified in the bottom row of Fig. 12. In CDM, there is a clear, monotonic correlation: high

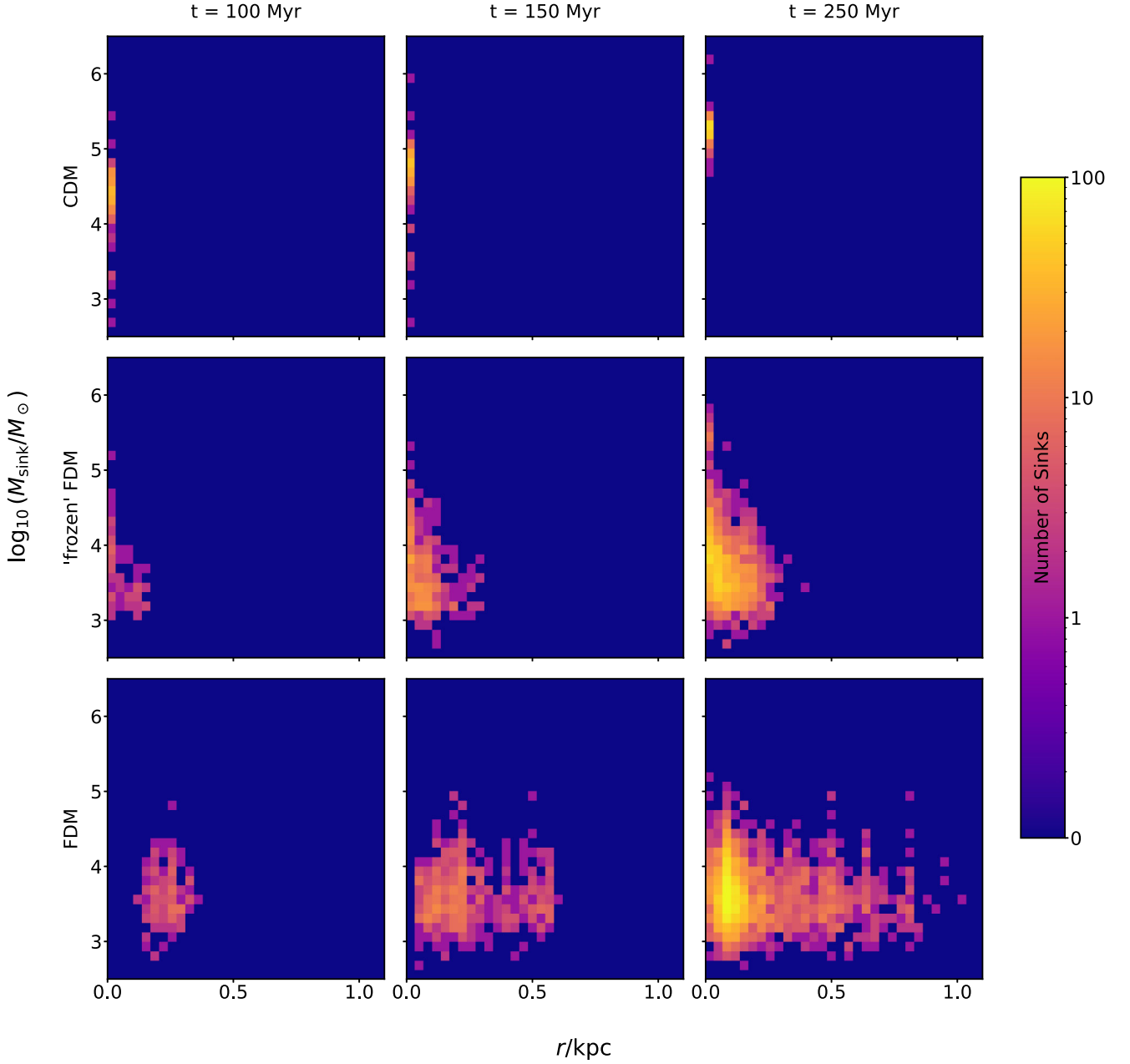


Figure 10. Time evolution of the sink particle population density in the mass–radius plane (M_{sink} versus r) for a halo of mass $M_h = 3 \times 10^9 M_\odot$. Rows display the results for CDM (top), ‘frozen’ FDM with $m_a = 1 \times 10^{-22}$ eV (middle), and the corresponding ‘dynamic’ FDM (bottom). The populations are shown at three characteristic epochs: shortly after first sink formation (100 Myr), intermediate (150 Myr), and late (250 Myr). The colour bar indicates the number of sink particles per bin on a logarithmic scale. In the CDM case, the population remains a highly concentrated small number of sinks at the halo centre, showing a vertical growth track as sinks continue to accrete mass in the deep central potential. In contrast, the FDM cases show a distinct broadening of the radial distribution. While sinks in the ‘frozen’ potential tend to remain tightly around the static gravitational centre over time, the fully dynamical FDM case exhibits significant migration of sinks to larger radii.

H_2 fractions only exist at high densities $\rho \gtrsim 10^9 M_\odot \text{ kpc}^{-3}$ ($n_{\text{H}} \gtrsim 10 \text{ cm}^{-3}$). In the ‘dynamic’ FDM case, the distribution of high H_2 concentration spreads much more broadly to the left, i.e. lower densities; high H_2 concentrations ($n_{\text{H}_2}/n_{\text{H}} \sim 10^{-2}$) persist at densities as low as $\rho \sim 10^4 M_\odot \text{ kpc}^{-3}$ ($n_{\text{H}} \sim 10^{-3} \text{ cm}^{-3}$).

The ‘frozen’ FDM case serves as a critical control. While it shows more extended H_2 than CDM due to its cored potential, it lacks the broad, low-density ‘tail’ seen in the ‘dynamic’ run. This confirms that the H_2 is not simply forming at larger radii due to

the core shape, but is being actively moved by the time-varying gravitational potential.

5.8.2 Implications for fragmentation

While we do not follow the gas down to higher densities in our simulations, this redistribution of cold gas has significant implications for the resulting stellar population. The thermodynamic state of primordial gas is a primary determinant of the fragmen-

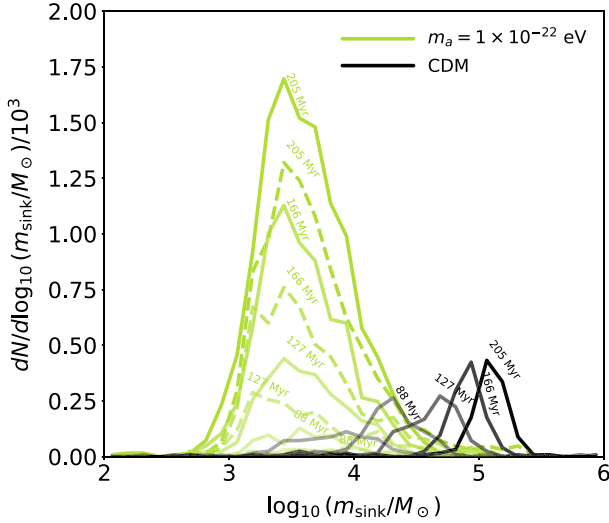


Figure 11. Evolution of the sink mass function for our $3 \times 10^9 M_\odot$ halo, showing the FDM ($m_a = 1 \times 10^{-22}$ eV), ‘frozen’ FDM (dashed), and CDM (black) cases over time. It should be noted that the sinks do not represent individual stars, but cold dense gas with the potential to form stars. The sink mass function should therefore not be overinterpreted. The CDM case shows a rapid increase in the mass of each sink particle, with a much larger average sink mass by the end of the simulation, while the number of particles stays roughly the same. In contrast, in the ‘dynamic’ FDM and ‘frozen’ FDM cases, the number of sinks increases with time, but their masses do not evolve, resulting in an abundant population of low-mass sinks.

tation scale (V. Bromm & R. B. Larson 2004; S. Glover 2013; R. S. Klessen & S. C. O. Glover 2023). By spreading the H_2 coolant over a wider volume, FDM dynamics create a more fragmented and turbulent environment. The presence of cold gas at lower densities suggests that the local Jeans mass, which scales as $M_J \propto T^{3/2} \rho^{-1/2}$, is reduced in the outskirts compared to a centrally concentrated CDM model. This potentially favours the formation of a more ‘bottom-heavy’ IMF or a more spatially dispersed cluster of stars compared to the single, massive central stellar concentration predicted by standard CDM scenarios (A. Stacy et al. 2011; P. C. Clark et al. 2011). This aligns with our observation in Section 5.7 that FDM sinks are scattered and prevented from accreting into high-mass objects.

6 DISCUSSION

Our simulations reveal that the wave-like nature of FDM does not merely provide a static cored potential, but introduces a complex dynamical environment that fundamentally reshapes the birthplaces of the first stars. While the monolithic $m_a \sim 10^{-22}$ eV model is increasingly constrained by observations, our results establish a crucial first step to characterize the interplay between wave-driven DM fluctuations and baryonic dynamics and chemistry. By treating this mass range as a physical upper bound for wave-driven effects, we have identified a new mechanism that suppresses gas accretion and redistributes key coolants. Below, we discuss the broader implications of these findings, their relevance to more complex DM cosmologies, and how these results compare to the wider cosmological landscape.

6.1 Scope and limitations

Our simulations demonstrate that in isolated haloes, the inherent dynamical behaviour of FDM suppresses the accumulation of cold dense gas compared to CDM. These results represent a foundational step in understanding gas collapse in ultra-light DM haloes. The simplifying assumptions made here (specifically a spherically symmetric halo with gas collapsing from rest) do not capture the more complex, anisotropic structure formation reproduced by cosmological simulations, where gas possesses initial kinetic energy and turbulence, and collapsed structures are rarely self-contained or spherically symmetric (e.g. P. Mocz et al. 2019). However, because the stochastic dynamics of the soliton are a universal feature of virialized FDM structures, we expect the qualitative conclusions regarding suppressed efficiency to be transferable to more comprehensively modelled, realistic FDM cosmologies. However, further work simulating these more realistic cosmological environments is required to confirm this result.

This study targets the low-mass range of haloes, 3×10^8 to $8 \times 10^9 M_\odot$, in which primordial stars form in FDM with $m_a = 10^{-22}$ to 10^{-21} eV considered here. However, these are still more massive than the $\sim 10^6 M_\odot$ minihaloes expected for realistic primordial star formation in Λ CDM. Such haloes are heavily suppressed in FDM for the m_a range that we consider in this paper. We would need to consider heavier axions, if we were to simulate $10^6 M_\odot$ haloes. However, such scenarios would require a different setup with much higher resolution, which would be computationally challenging in this setup, or a much smaller box size. Furthermore, these first star formation simulations must eventually include several feedback mechanisms unique to molecular cooling haloes, such as Lyman–Werner feedback (Z. Haiman, M. J. Rees & A. Loeb 1997), the suppression of star formation by relative motion between DM and gas (A. T. P. Schauer et al. 2021), and internal turbulence (A. Stacy et al. 2011), which we leave for future work.

We have identified a region in the M_h – m_a parameter space where FDM dynamics create signatures distinct from CDM. However, a finer sampling in m_a , particularly in the regime 1×10^{-22} eV $\lesssim m_a \lesssim 3 \times 10^{-22}$ eV where the most interesting deviations from CDM were found, would allow for more reliable calibrations of star formation models. This change in gas distribution is especially evident in the transition from growing to shrinking sink half-mass radii observed in our simulations.

6.2 Fragmentation, chemical stirring, and the Population III IMF

A unique finding of this work is the chemical stirring effect (Section 5.8), where FDM fluctuations transport H_2 -rich gas into low-density regions. While DM dynamics suppress the total gas accumulation, this redistribution suggests that star formation may be more spatially dispersed than in the monolithic CDM case. The presence of extra coolant far from the halo centre may have significant implications for the resulting stellar Population III IMF.

In CDM, the stable central potential favours the formation of a few massive sinks via sustained accretion. In contrast, the scattering and rotational support we observe in FDM results in a more fragmented, lower mass sink population. It is reasonable to anticipate that such a significant modification to the sink population will be reflected in the characteristics and distribution of the first stars.

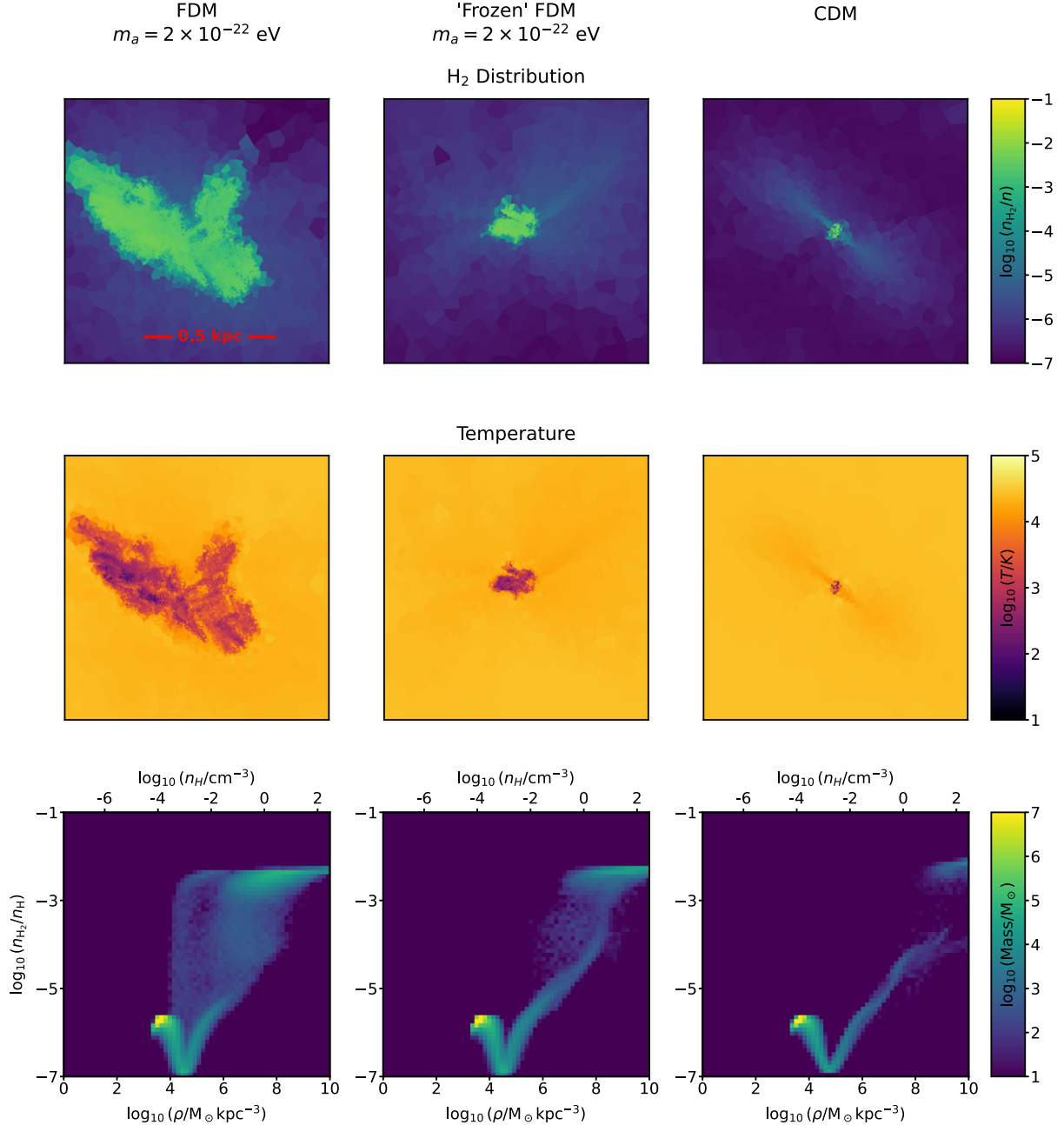


Figure 12. H_2 properties at $t \approx t_* + 190 \text{ Myr}$ (synchronized with the final row of Fig. 5), shown for a slice through a $M_h = 8 \times 10^8 M_\odot$ halo centre for FDM with $m_a = 2 \times 10^{-22} \text{ eV}$ (left), the corresponding ‘frozen’ FDM case (centre), and the comparison CDM case (right). We demonstrate the distribution of H_2 (top), the gas temperature of the same slice (middle), and the H_2 concentration–density phase space (bottom). It is clear that in the FDM case, gas with a high concentration of H_2 ($n_{\text{H}_2}/n_{\text{H}} > 10^{-3}$) is liberated much further from the halo centre than in either the ‘frozen’ FDM or the CDM case. This is evident from the bottom panels, where in the FDM case (left), concentrations of $n_{\text{H}_2}/n_{\text{H}} \gtrsim 10^{-3}$ extends to much lower gas densities than either the ‘frozen’ FDM or CDM cases.

While we can only speculate here about the effects of FDM dynamics on the true Pop III IMF, increased levels of turbulence have been shown to promote fragmentation and lead to smaller stellar masses (P. C. Clark et al. 2011; R. S. Klessen & S. C. O. Glover 2023); the wave-driven ‘stirring’ in FDM may similarly shift the Pop III IMF toward lower masses. Additionally, a broader region of gas enriched with H_2 may further aid cooling and fragmentation with implications for the Pop III IMF. However, further work is needed to follow the gas down to higher densities and

smaller scales to isolate individual stars before robust conclusions can be made about any effect on the Pop III IMF.

6.3 Comparison with baryonic feedback and different dark matter models

To contextualize the effects of the FDM dynamics demonstrated in this work, it is critical to evaluate how alternative solutions to the small-scale challenges of ΛCDM reshape the baryonic

components, and whether a degeneracy exists between the effects of the FDM dynamics and baryonic feedback mechanisms.

Cosmologists have looked to resolve these challenges with much success by invoking baryonic physics, such as stellar feedback from bursty star formation, supernovae explosions, and active galactic nuclei activity (S. Mashchenko, J. Wadsley & H. M. P. Couchman 2008; T. Okamoto, L. Gao & T. Theuns 2008; F. Governato et al. 2012; A. Pontzen & F. Governato 2012; A. R. Wetzel et al. 2016; S. Garrison-Kimmel et al. 2019; L. V. Sales et al. 2022; A. Schneider et al. 2025). It remains unclear how the conclusions of this work would be affected with a more detailed treatment of these feedback mechanisms, and we leave this open question to future work.

Self-interacting dark matter (SIDM) can address some of the small-scale challenges by creating flattened density cores through elastic particle scattering (D. N. Spergel & P. J. Steinhardt 2000; M. Rocha et al. 2013; M. Kaplinghat, S. Tulin & H.-B. Yu 2016; S. Tulin & H.-B. Yu 2018; S. Adhikari et al. 2025). While both FDM and SIDM result in a suppression of the central gravitational cusp, they predict different relationships between halo mass and core properties, such as central density and core radius. Since our results demonstrate that gas accumulation and fragmentation are highly sensitive to the specific geometry and depth of the central potential, it remains a compelling open question how the distinct core profiles of SIDM would affect the timing and scale of primordial star formation. A detailed comparison between the wave-driven dynamics of FDM and the scattering-induced profiles of SIDM would be essential for identifying unique baryonic signatures that could distinguish between these two DM candidates.

Self-interacting fuzzy dark matter blends the ultra-light, wave-like nature of FDM with the collisional properties of SIDM, potentially alleviating the core-density tensions found in single-field models (P. Mocz et al. 2023; M. Indjin et al. 2024; C. Capanelli, W. Hu & E. McDonough 2025; M. Indjin et al. 2026b). These scalar fields generically have quartic self-interactions that can modify the evolution of density perturbations, leading to a more effective flattening of central density profiles. While exploring the full parameter space of self-interactions and multifield DM remains a broader unexplored area, isolating the dynamics of the non-self-interacting, single-field FDM limit here establishes the foundations required to systematically evaluate these next-generation simulations.

Many of the tight constraints on m_a are ultimately predicated on the assumption that a single axion-like species constitutes ~ 100 per cent of the DM. MDM models, which combine CDM and FDM or multiple FDM fields of different masses (M. Gosenca et al. 2023; H. N. Luu et al. 2024; T. Dome et al. 2025), provide a promising alternative pathway to reconcile these observations. As shown by T. Dome et al. (2025), axion depletion in haloes below the FDM Jeans mass may effectively insulate dwarf galaxies from dynamical heating, while still allowing for a globally significant 10^{-22} eV component. The findings of this work therefore provide the necessary calibration, reference data, and essential single-field baseline required to interpret these more observationally viable multicomponent cosmologies. Exactly how our conclusions would be affected in such MDM haloes remains unclear and is a question we hope to address in future work.

6.4 Implications for *JWST* and 21-cm cosmology

The suppressed star formation efficiency has direct consequences for large-scale observable signals. A delay or reduction in early

star formation due to FDM dynamics naturally leads to a delay in the onset of both Cosmic Dawn and the Epoch of Reionization (EoR), thus shortening their duration relative to what is predicted in Λ CDM (O. Nebrin et al. 2019; D. Jones et al. 2021; S. Liu et al. 2025; J. Dhandha et al. 2025b). As a result, the cosmic thermal and ionization histories are altered, impacting the large-scale cosmological 21-cm signal of neutral hydrogen. Specifically, the modifications to the Population III IMF driven by FDM dynamics are expected to propagate into the X-ray heating of the intergalactic medium (N. S. Sartorio et al. 2023). This suggests a potentially stronger impact of FDM on the 21-cm signal than previously predicted, as the sensitivity of the 21-cm global signal and power spectrum to the Pop III IMF has been shown to be significant in the CDM context (T. Gessey-Jones et al. 2025). Consequently, future 21-cm experiments may provide robust avenues to constrain the axion mass to within 10 per cent around 10^{-21} eV (S. Liu et al. 2025).

The recent discovery by *JWST* of an unexpectedly high number of massive, bright galaxies at $z > 10$ has sparked significant debate regarding the efficiency of early star formation. While the general delay in star formation found in FDM models might initially appear to exacerbate this tension, it is important to note that current observations primarily constrain relatively massive haloes ($M_h > 10^9 M_\odot$). At these scales, our simulations show a transition toward CDM-like behaviour, suggesting that the most luminous sources may be less affected by wave-like dynamics. However, the mass- and time-dependent star formation efficiency we have uncovered remains a critical variable; as suggested by J. Dhandha et al. (2025a), a flexible efficiency that scales with halo mass can help reconcile theoretical predictions with *JWST* counts. In the low-mass regime where FDM dynamics are most potent, our predicted suppression remains an indicator for the faint end of the UVLF under FDM conditions.

7 CONCLUSIONS

In this study, we performed high-resolution hydrodynamical simulations of FDM using AXIREPO to investigate the impact of wave-like DM dynamics on the collapse and cooling of primordial gas. By comparing full dynamical simulations to ‘frozen FDM’ and CDM counterparts, we isolated the role of solitonic fluctuations in regulating star formation. Our simulations target haloes in the mass range 3×10^8 to $8 \times 10^9 M_\odot$, a regime where the processes of halo formation, gas collapse, and star formation are strongly affected by FDM physics for axion masses of $m_a = 1 \times 10^{-22}$ to 7×10^{-22} eV. We treat this range not as a claim for a monolithic DM candidate, but as a critical numerical benchmark to establish the maximal possible impact of wave DM in these haloes.

We find that the stochastic random walk and mass fluctuations of the solitonic core act as a continuous source of gravitational perturbation. This leads to a substantial delay in the onset of first sink formation (the proxy for dense star-forming gas) and a significant reduction in total accumulated sink mass. While the delay in structure formation due to the FDM power spectrum cut-off is well documented (H.-Y. Schive et al. 2014a; S. May & V. Springel 2021, 2023), we show for the first time through direct simulation that the internal dynamics of the halo provide a secondary, independent suppression mechanism that further delays and suppresses the onset of Cosmic Dawn. By simulating the regime where FDM dynamics are most prominent, we establish an upper bound on the potential impact of FDM on the primordial environment. If we had found that the gas

physics and star formation remain relatively unaffected in this extreme limiting case, it would provide a strong indication that such effects are negligible in more conservative or higher mass models.

Furthermore, our work reveals that FDM dynamics induce significant specific angular momentum in the gas within the soliton radius. For some parameter combinations at the low-mass end of the m_a and M_h range explored, this acquired angular momentum reaches values comparable to the critical threshold ($j \gtrsim j_{\text{crit}}$) required for rotational stabilization out to the soliton radius. This provides a physical barrier to collapse that is entirely absent in the CDM and ‘frozen’ FDM cases, demonstrating that soliton dynamics, rather than just the core’s geometry, dictate the spatial scale of gas accumulation.

Our results demonstrate that the formation of the first sink is delayed in FDM compared to CDM counterparts, with the magnitude of this delay scaling inversely with both halo mass and axion mass. By comparing our ‘frozen’ and ‘dynamic’ models, we find that the flattened potential gradient of the cored FDM profile in the ‘frozen’ case is sufficient to explain this delay in most cases. However, the full wave-like fluctuations in the ‘dynamic’ models can cause an additional delay in the onset of star formation at the low-mass end of the m_a and M_h range examined in this work. In these cases, the FDM simulations experience a greater delay in first sink formation than their ‘frozen’ counterparts, confirming that soliton fluctuations provide a secondary, dynamical barrier that significantly delays gas collapse beyond what a static cored-potential would predict.

Within our simulations, we also see a dramatic shift in the distribution and growth of our sink particles, which in our simulations serve as proxies for dense star-forming regions. In CDM, sinks remain centrally concentrated and accrete mass efficiently. In our FDM haloes (both ‘frozen’ and ‘dynamic’), sinks form at larger radii (near the soliton radius r_c). However, in the ‘dynamic’ FDM case, these sinks are progressively scattered out to higher radii than in the ‘frozen’ comparison case.

We observe that FDM fluctuations extend the distribution of molecular hydrogen (H_2) and cold gas into lower density regions. This chemical stirring induces further cooling over a broader volume than is seen in CDM haloes. While the total gas collapse is suppressed, this dispersed cooling indicates a shift from monolithic central star formation to more fragmented, spatially diffuse clusters of dense, cold gas. This, coupled with our findings regarding the distribution and growth of sinks, hints that low-mass FDM haloes may naturally induce greater fragmentation and hence shift the Population III IMF toward lower masses compared to the massive stars predicted in CDM.

Our results provide an additional physical mechanism, beyond the large-scale power spectrum cut-off, leading to a delayed EoR in pure FDM cosmologies, and motivates further study of these effects in more complex MDM models. We suggest that the internal inefficiency of gas collapse, driven by both solitonic geometry and wave dynamics, further slows the contribution of the smallest FDM haloes to the ionizing budget. Critically, because this dynamical suppression is strongly mass-dependent and diminishes in more massive systems, our model suggests a high degree of variance in star formation efficiency across the halo mass function. This mass-dependent scaling offers a potential pathway to reconcile FDM with recent *JWST* observations of massive high-redshift galaxies (J. Dhandha et al. 2025a); while FDM strongly inhibits star formation in the smallest dwarf precursors, it allows more massive haloes to assemble and form stars with efficiencies

approaching those of Λ CDM, effectively shifting the balance of early light production toward the most massive systems.

Finally, our work demonstrates the necessity of incorporating accurate DM dynamics when deriving constraints from baryonic observables. This dynamical framework is essential for future investigations into multifield or self-interacting FDM models, as we move toward a more complete description of DM that better reflects the constraints on single-field FDM models.

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DATA AVAILABILITY

Snapshot data and post-processed results are made available upon reasonable request.

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APPENDIX A: SINK FORMATION AND EVOLUTION IN RADIUS-MASS SPACE

In this appendix, we provide a visualization of the sink particle lifecycle within our simulated haloes. Fig. A1 shows the distribution of a random sample of 16 sink particles in the radius-mass plane, specifically tracking their initial formation (shown in the circles) and their final state at the end of the simulation (shown as the end of the arrows). These data were taken from our $3 \times 10^9 M_\odot$ halo.

The figure highlights the distinct growth modes discussed in Section 5.7. In CDM, sinks are primarily born near the centre of the deep gravitational cusp and rapidly accrete gas. In contrast, the FDM haloes show sink formation across a wider range of radii, coinciding with the interference-driven fluctuations of the solitonic core. The evolutionary tracks (arrows) in FDM show significantly less mass growth, and a general movement to larger radii (gravitational heating) confirming that wave-driven turbulence and angular momentum support limit the efficiency of gas accretion on to individual pre-stellar cores. In contrast, while the ‘frozen’ FDM is able to form sinks at larger radii than CDM, these sinks generally stay closer to the centre of the halo.

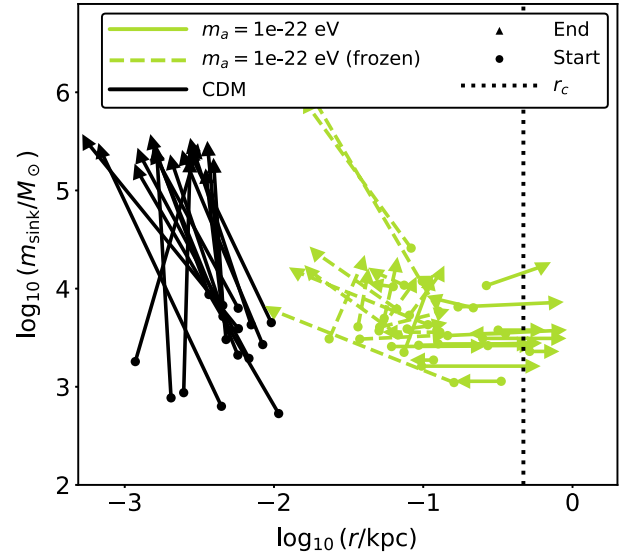


Figure A1. The distribution of sink particles in radius-mass space for representative CDM (black) and FDM (green) haloes. Circles denote the properties of sinks at the moment of formation, while the arrow ends indicate their final mass and radial position. Arrows indicate the evolutionary path of individual sinks, highlighting mass accretion and radial migration. In the CDM case, sinks quickly migrate toward the centre and grow into massive objects, whereas in the FDM cases, sinks remain more distributed with lower final masses, particularly in the ‘dynamic’ FDM (solid arrows) where a general migration outwards is seen.

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