Dark Matter through the Axion Portal

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Motivated by the galactic positron excess seen by PAMELA and ATIC/PPB-BETS, we propose that dark matter is a TeV-scale particle that annihilates into a pseudoscalar "axion." The positron excess and the absence of an anti-proton or gamma ray excess constrain the axion mass and branching ratios. In the simplest realization, the axion is associated with a Peccei-Quinn symmetry, in which case it has a mass around 360 – 800 MeV and decays into muons. We present a simple and predictive supersymmetric model implementing this scenario, where both the Higgsino and dark matter obtain masses from the same source of TeV-scale spontaneous symmetry breaking.

Introduction. Evidence for dark matter (DM) is by now overwhelming [1]. While the precise nature and origin of DM is unknown, thermal freezeout of a weakly interacting massive particle (WIMP) is a successful paradigm that arises in many theories beyond the standard model. If this is correct, then specific DM properties can be probed through direct and indirect detection experiments, and pieces of the dark sector might even be produced at the Large Hadron Collider (LHC).

Recent indirect detection results may offer important insights into the dark sector. The latest PAMELA data [2] is strongly suggestive of a new source of galactic positrons, bolstering the HEAT [3] and AMS-01 [4] anomalies. An intriguing interpretation of the PAMELA excess is DM annihilation [5, 6, 7, 8]. The rate and energy spectrum of the PAMELA positrons are consistent [7] with the electron/positron excesses seen in the balloon experiments ATIC [9] and PPB-BETS [10], and the spectral cutoff in these experiments points to a DM mass in the TeV range.

There are, however, two puzzling features in the PAMELA data. First, the positron excess is not accompanied by an anti-proton excess, which strongly constrains the hadronic annihilation modes of the DM [7]. Unless DM is as heavy as 10 TeV, the PAMELA data disfavors DM annihilation into quarks, Ws, Zs, or Higgses. Second, the required annihilation cross section in the galactic halo is orders of magnitude larger than the thermal relic expectation. Therefore, any DM interpretation of the data must explain both the large annihilation rate and the large annihilation fraction into leptons.

In this letter, we propose that DM is a TeV-scale particle that dominantly annihilates into a pseudoscalar "axion" a. The axion mass lies above the electron or muon threshold, so that a dominantly decays into leptons with suppressed photonic and hadronic decay modes. In this scenario, the electron/muon decay channels would account for the PAMELA excess, and the photon/pion/tau decay channels would be constrained by gamma ray telescopes like HESS [11] and FERMI [12].

In a typical realization of the scenario, DM is a fermion. In this case, the dominant annihilation channel is not to $2a$, but to a and a real scalar s. If the scalar dominantly decays as $s \to aa$, each DM annihilation will contain three axions, yielding a distinctive semi-hard galactic positron spectrum. The existence of a light scalar s is also crucial to enhance the DM galactic annihilation rate through nonperturbative effects.

This DM scenario can arise in any theory where the DM mass is generated from spontaneous breaking of a global $U(1)_X$ symmetry under which leptons have axial charges. In two Higgs doublet models, it is natural to identify $U(1)_X$ with a Peccei-Quinn (PQ) symmetry $U(1)_{\rm PQ}$ [13], in which case our axion is a heavier variant of the DFSZ axion [14]. In this realization, the axion must decay dominantly into muons to evade constraints from low energy and astrophysical experiments.

The identification of $U(1)_X$ with $U(1)_{\rm PQ}$ is particularly well-motived in supersymmetric (SUSY) theories, since then both the Higgsino and DM can obtain masses from the same source of TeV-scale spontaneous $U(1)_{\rm PQ}$ breaking. As we will see, the resulting SUSY model is extremely simple and predictive, and leads to interesting phenomenology at the LHC. We thus mainly focus on this model, although other possibilities are also discussed.

Basic Setup. To understand the main features of the scenario, we first isolate the fields responsible for the dominant DM phenomenology and study their dynamics. A complete SUSY model will be described later.

Our starting point is a global $U(1)_X$ symmetry that is broken by the vacuum expectation value (vev) of a complex scalar S:

$$S = \left(f_a + \frac{s}{\sqrt{2}} \right) e^{ia/\sqrt{2} f_a}, \tag{1}$$

where a is the "axion," s is a real scalar, and f_a is the axion decay constant. A vector-like fermion DM ψ/ψ^c obtains a mass from

$$\mathcal{L} = -\xi S \psi \psi^c + \text{h.c.}, \qquad m_{\rm DM} = \xi f_a, \tag{2}$$

where ψ/ψ^c is a standard model (SM) singlet. The stability of DM can be ensured by a vector-like symmetry acting on ψ/ψ^c, which could be a remnant of $U(1)_X$ [31].

In order for a to decay into leptons, SM leptons must have nontrivial $U(1)_X$ charges. For example, in a two Higgs doublet model, a coupling of the form

$$\mathcal{L} = f(S) h_u h_d + \text{h.c.}, \tag{3}$$

can force $h_u h_d$, and hence the quarks and leptons, to carry nontrivial $U(1)_X$ charges. The $U(1)_X$ is then a PQ symmetry, and we focus on this case until the end of this letter. If the coupling of Eq. (3) is sufficiently small, it does not drastically affect the DM phenomenology.

Unlike the ordinary axion [15], the mass of a cannot come only from pion mixing, since it would then be too light to decay into leptons. We therefore need some explicit breaking of $U(1)_X$. Also, the potential that generates the S vev will also give a mass to s. Both effects can be described phenomenologically by the mass terms

$$\mathcal{L} = -\frac{1}{2} m_a^2 a^2 - \frac{1}{2} m_s^2 s^2, \tag{4}$$

and we assume the hierarchy $m_a \ll m_s \ll m_{\rm DM}$. This condition is naturally satisfied in the explicit SUSY model considered later.

The scalar field s decays into aa through the operator

$$\mathcal{L} = \frac{1}{\sqrt{2} f_a} s (\partial_\mu a)^2, \tag{5}$$

arising from the S kinetic term. This is typically the dominant decay channel for s.

Thermal Freezeout. In the above setup, the DM has three major annihilation modes

$$\psi\bar{\psi} \to ss, \qquad \psi\bar{\psi} \to aa, \qquad \psi\bar{\psi} \to sa. \tag{6}$$

The first two modes do not have an s-wave channel and are suppressed in the $v \to 0$ limit. Annihilation at thermal freezeout is therefore dominated by the third mode.

In the limit $m_s, m_a \ll m_{\rm DM}$, the $v \to 0$ annihilation cross section is

$$\langle \sigma v \rangle_{\psi\bar{\psi} \to sa} = \frac{m_{\rm DM}^2}{64\pi f_a^4} + O(v^2). \tag{7}$$

A standard thermal relic abundance calculation implies

$$\langle \sigma v \rangle = \frac{1}{2} \langle \sigma v \rangle_{\psi\bar{\psi} \to sa} \simeq 3 \times 10^{-26} \text{cm}^3/\text{s}, \tag{8}$$

so once $m_{\rm DM}$ is constrained by future ATIC data, then f_a is completely determined. As a fiducial value, $m_{\rm DM} \sim$ 1 TeV implies $f_a \sim 1$ TeV, and so $\xi \sim 1$.

Halo Annihilation. The cross section of Eq. (8) is too small to account for the observed PAMELA excess. For $m_{\rm DM} \sim 1$ TeV, the required boost factor is [6, 7]

$$\langle \sigma v \rangle_{\rm PAMELA} \simeq 10^3 \langle \sigma v \rangle, \tag{9}$$

although the precise value is subject to a factor of a few uncertainty. Such a large boost factor is difficult to explain astrophysically [16].

However, the halo annihilation rate can be enhanced by nonperturbative effects associated with the light state s, with $m_s \ll m_{\rm DM}$ [8, 17]. The relevant effects are the Sommerfeld enhancement and the formation of DM boundstates (WIMPoniums). These boost the signal by

$$B \simeq c \frac{\alpha_\xi m_{\rm DM}}{m_s}, \qquad \alpha_\xi = \frac{\xi^2}{4\pi}, \tag{10}$$

where c is a coefficient which can be as large as $(m_s/\alpha_\xi m_{\rm DM})^{3/2} \alpha_\xi^2 / v_{\rm halo}^2$ if m_s takes values that allow (near) zero-energy boundstates. Here, $v_{\rm halo} \sim 10^{-3}$. As we will see, our explicit model has $m_s \approx O(1 - 10 \text{ GeV})$. Combined with a moderate astrophysical boost factor, Eq. (10) can then account for Eq. (9). The condition for the effects being operative in the halo, but not at freezeout, is $v_{\rm halo} \lesssim \alpha_\xi \lesssim v_{\rm freezeout}$, where $v_{\rm freezeout} \sim 0.4$ [32].

The DM annihilation (or para-WIMPonium decay) product is sa. After the decay $s \to aa$, this yields three axions per DM annihilation. There is also an annihilation channel into $t\bar{t}$ through s-channel a exchange, but its branching fraction is only of $O(1\%)$. This level of hadronic activity is consistent with the PAMELA data.

Since DM does not annihilate directly into leptons, the positron injection spectrum is different from [7] and closer to [19]. If a DM annihilation product comes from a (scalar) cascade decay with large mass hierarchies involving n steps, then its energy spectrum is proportional to $\{\ln(m_{\rm DM}/E)\}^{n-1}$, where $n = 0$ for direct annihilation. The lepton spectrum per DM annihilation is then

$$\frac{dN_\ell}{dE} = \frac{2}{m_{\rm DM}} \left(2 \log \frac{m_{\rm DM}}{E} + 1 \right) \quad (E < m_{\rm DM}). \tag{11}$$

This semi-hard lepton spectrum is an unambiguous prediction of our setup. Since our axion will primarily decay into muons, Eq. (11) must be convoluted with muon decays to give the final positron spectrum.

Axion Decays. To account for the observed positron excess, the axion must dominantly decay into leptons. As we will see, there are strong bounds on the photon flux from DM annihilations and this constrains $a \to \gamma\gamma$ to be less than $\approx 1\%$. Since $\pi^0 \to \gamma\gamma$, the decay into neutral pions must also be suppressed to the 5% level. This eliminates the possibility $a \to \tau^+\tau^-$, since tau decay leads to an $O(1)$ fraction of π^0s.

Compared to the QCD axion, we have an extra m_a^2 parameter which affects axion-pion mixing. In terms of the mixing angle for the QCD axion $\bar{\theta}_{a\pi^0} \sim f_\pi/f_a$, the new mixing angle is $\theta_{a\pi^0} = \bar{\theta}_{a\pi^0} m_\pi^2/(m_\pi^2 - m_a^2)$. Close to the π^0 threshold there is resonant enhancement, but for $m_a^2 \gg m_\pi^2$, there is a m_π^2/m_a^2 suppression. A similar enhancement also arises for $m_a^2 \simeq m_\eta^2$.

For a generic axion, the partial widths to leptons and photons are

$$\Gamma(a \to \ell^+\ell^-) \ = \ c_\ell^2 \frac{m_a}{16\pi}\frac{m_\ell^2}{f_a^2}, \qquad (12)$$

$$\Gamma(a \to \gamma\gamma) \ = \ c_\gamma^2 \frac{\alpha_{\rm EM}^2}{128\pi^3}\frac{m_a^3}{f_a^2}. \qquad (13)$$

In our case, $c_\ell = \sin^2\beta$, where $\tan\beta \equiv \langle h_u\rangle/\langle h_d\rangle$, and c_γ depends on m_a through $\theta_{a\pi^0}$ but is typically $O(1)$. Except when m_a is very far from a lepton mass threshold, the photon branching fraction is less than $O(1\%)$.

The partial width to pions is more complicated. Direct decays $a \to \pi\pi$ are suppressed by CP invariance, and radiative decays $a \to \pi\pi\gamma$ are suppressed by $\alpha_{\rm EM}/4\pi$. The first dangerous channel is $a \to \pi\pi\pi$ which arises from axion-pion mixing. We estimate the partial width as

$$\Gamma(a \to \pi\pi\pi) \sim \frac{1}{128\pi^3}\frac{m_a^5}{f_\pi^4}\left(\frac{f_\pi}{f_a}\frac{m_\pi^2}{m_a^2}\right)^2, \qquad (14)$$

where the combination in parentheses is the approximate axion-pion mixing angle when $m_a^2 \gg m_\pi^2$. Parametrically, the $a \to \pi\pi\pi$ mode is two orders of magnitude suppressed compared to the $a \to \mu^+\mu^-$ mode. As m_a approaches the $\rho\pi$ threshold, $a \to \pi\pi\pi$ is enhanced by m_ρ^2/Γ_ρ^2. Also, $a \to \eta\pi\pi$ decay becomes important around the same mass scale, so we estimate the total axion to π^0 branching ratio to be safe for $m_a \lesssim 800$ MeV.

Axion Constraints. The bounds on heavy axions are different from ordinary axions. For the range $2m_e < m_a < 2m_\mu$, a beam-dump experiment at CERN [20] looked for the decay $a \to e^+e^-$, and definitively rule out axion decay constants up to $f_a \sim 10$ TeV.

In the region $2m_\mu < m_a < m_K - m_\pi$, our axion decays into $\mu^+\mu^-$ with $c\tau_a \simeq O(1-10\ \mu{\rm m})$, and measurements of rare kaon decays $K \to \pi\mu^+\mu^-$ constrain the branching ratio $K \to \pi a$. The estimated branching ratio is ${\rm Br}\,(K^+ \to \pi^+ a) \gtrsim 1\times10^{-4}(v_{\rm EW}/f_a)^2$ [21], and the measured rate ${\rm Br}\,(K^+ \to \pi^+\mu^+\mu^-) \simeq 1\times 10^{-7}$ [22] is consistent with SM expectations. Therefore, this region seems to be excluded for $f_a \sim 1$ TeV, especially considering that the dimuon invariant mass distribution would be peaked at m_a for the axion decay.

For $m_K - m_\pi < m_a \lesssim 800$ MeV, there are bounds from $\Upsilon \to \gamma a$ where a decays promptly into $\mu^+\mu^-$ [23]. These searches, however, impose no constraint for the region of interest $f_a \sim 1$ TeV. Note that many of the J/Ψ and Υ decay experiments looking for the γa mode do not apply here because our axion decays promptly.

In summary, the allowed region for our axion is

$$m_K - m_\pi < m_a \lesssim 800 \text{ MeV}, \qquad (15)$$

and the dominant decay channel is $a \to \mu^+\mu^-$. For axions as heavy as $m_K - m_\pi$, astrophysical bounds are irrelevant.

Gamma Ray Spectrum. As already mentioned, the axion typically has a non-zero branching fraction into photons (or π^0s), and there are important bounds from gamma ray experiments. Since the photon spectrum is semi-hard, the strongest bounds come from atmospheric Cherenkov telescopes. The expected photon spectrum also overlaps with the energy range of FERMI.

A model-independent constraint comes from a HESS study of the Sagittarius dwarf galaxy [24]. They put an upper bound on the integrated gamma ray flux for $E_\gamma > 250$ GeV of $\Phi_\gamma < 3.6\times10^{-12}\ {\rm cm}^{-2}\,{\rm s}^{-1}$. From a given photon spectrum, this can be translated into a bound on the DM annihilation cross section

$$\langle\sigma v\rangle < \frac{4\pi\Phi_\gamma m_{\rm DM}^2}{\bar{J}\Delta\Omega}\left(\int_{250\ {\rm GeV}}^{m_{\rm DM}}\frac{dN_\gamma}{dE}\,dE\right)^{-1}, \qquad (16)$$

where $\bar{J} \simeq 2.2\times10^{24}\ {\rm GeV}^2\,{\rm cm}^{-5}$ is the Sagittarius line-of-sight-integrated squared DM density assuming an NFW profile, and $\Delta\Omega = 2\times10^{-5}$ is the HESS solid angle integration region.

Since the annihilation cross section is known from Eq. (9), we can translate Eq. (16) into a bound on the branching fraction to photons. For direct $a \to \gamma\gamma$ decays, the energy spectrum is proportional to Eq. (11), $dN_\gamma/dE \simeq {\rm Br}(a \to \gamma\gamma)\,dN_\ell/dE$, and using the fiducial $m_{\rm DM} = 1$ TeV, we obtain the bound

$${\rm Br}(a \to \gamma\gamma) \lesssim 1\% \quad (m_{\rm DM} = 1\ {\rm TeV}). \qquad (17)$$

The bound on axion decays into pions can be derived similarly. Assuming $a \to \pi^0\pi^+\pi^-$, we obtain

$${\rm Br}(a \to \pi\pi\pi) \lesssim 5\% \quad (m_{\rm DM} = 1\ {\rm TeV}). \qquad (18)$$

The decay $a \to \tau^+\tau^-$ leads to a bound on $\langle\sigma v\rangle$ an order of magnitude stronger than $\langle\sigma v\rangle_{\rm PAMELA}$ [33].

Supersymmetric Model. In a SUSY context, it is natural to assume that the vector-like DM mass is related to the vector-like Higgsino mass μ_H. In fact, the simple superpotential

$$W = \xi S\Psi\Psi^c + \lambda S H_u H_d, \qquad (19)$$

together with the soft SUSY breaking terms

$$\mathcal{L}_{\rm soft} = -\xi A_\xi S\Psi\Psi^c - \lambda A_\lambda S H_u H_d - m_S^2 S^\dagger S + \cdots, \qquad (20)$$

have all the required ingredients, except for the origin of the axion mass, which we leave unspecified (can simply be a small κS^3 term in the superpotential). Without the ξ terms, this model is the PQ-symmetric limit of the NMSSM, and is sometimes referred to as PQ-SUSY [26, 27]. The vevs for S, H_u and H_d can be generated in a stable vacuum, giving $m_{\rm DM}/\mu_H = \xi/\lambda$.

For $\lambda \ll 1$ and $|m_S^2| \ll \lambda^2 v_{\rm EW}^2$, where $v_{\rm EW} \simeq 174$ GeV, the dominant phenomenology is determined essentially by five parameters

$$\{m_{\rm DM}, \lambda, \tan\beta, m_S^2, m_a\}, \qquad (21)$$

$m_{\rm DM}$	λ	$\tan\beta$	m_S^2	f_a	μ_H	A_λ	$m_{H_u}^2$	$m_{H_d}^2$	m_s	τ_s	$\mathrm{Br}(s \to f\bar{f})$	$m_{\tilde{s}}$	$m_{3/2}$	$\tau_{\tilde{s}}$	m_a	τ_a	$\sigma_{\rm SI}$ [cm^2]
1000	0.25	2.0	-6.8^2	1100	270	650	110^2	530^2	34	$4 \cdot 10^{-21}$	$f = b : 3\%$	5.5	10 eV	$2 \cdot 10^{-5}$	0.7	$8 \cdot 10^{-15}$	$3 \cdot 10^{-43}$
1200	0.10	4.0	-6.3^2	1200	120	430	-69^2	440^2	5.6	$1 \cdot 10^{-18}$	$f = \tau : 5\%$	1.2	5 eV	0.02	0.4	$1 \cdot 10^{-14}$	$4 \cdot 10^{-43}$

TABLE I: Two sample spectra in the SUSY model. $m_{\rm DM}$, λ, $\tan\beta$, m_S^2, $m_{3/2}$, and m_a are inputs, and the rest are outputs. All the masses are in GeV (except where indicated), and the lifetimes are in seconds. $\sigma_{\rm SI}$ is a spin-independent DM-nucleon cross section. $m_{\tilde{s}}$ is calculated assuming decoupling gauginos.

with m_S^2 and m_a affecting only scalar mixing and axion decay, respectively. All other parameters are either determined by thermal relic calculations, electroweak symmetry breaking, or are secondary to the phenomenology relevant here. The model is thus extremely predictive in this region, and we present sample spectra in Table I.

The present SUSY model introduces important additions to the minimal structure described before. First, the mass of s is no longer a free parameter and is fixed by $m_s \simeq \lambda\, v_{\rm EW} \sin(2\beta)$. Second, we have an additional light state $\tilde{s}$, the fermion component of S, whose mass is $m_{\tilde{s}} \simeq O(\lambda^2 v_{\rm EW}^2 / m_{\rm SUSY})$, where $m_{\rm SUSY}$ is μ_H or a gaugino mass.

The existence of S states lighter than the electroweak scale is consistent with the experimental data. These states mix with $H_{u,d}$ states with mixing angles of $O(v_{\rm EW}/f_a)$, and constraints from LEP are satisfied for $f_a \gtrsim 1$ TeV. Considering $\mu_H = \lambda f_a$ ($\simeq A_\lambda \sin(2\beta)/2$ from potential minimization), $f_a \sim 1$ TeV implies $\lambda \approx O(0.1)$, which satisfies the bound on charginos.

Small values of λ allow $m_s \approx O(1-10~\text{GeV})$, as needed for the halo annihilation enhancement. Small λ also ensures that DM annihilates mainly into S states and not $H_{u,d}$ states, which would give more hadronic activity than is allowed by PAMELA. To suppress additional hadronic/photonic activity from s decays, the branching fraction of s into quarks and taus can be made smaller than $O(10\%)$. The $s \to \tilde{s}\tilde{s}$ mode is subdominant, and annihilations of DM into $\tilde{s}\tilde{s}$ are velocity suppressed.

Since s is light, it mediates a large DM-nucleon cross section, leading to tension with the direct detection bound [28]. There are two ways this bound can be satisfied. One is to take m_s to be a few tens of GeV. In this case the halo annihilation enhancement occurs through (near) zero-energy boundstates. The other is to suppress the s-h_d mixing by taking appropriate values of $|m_S^2| \approx 0.1\lambda^2 v_{\rm EW}^2$. In this case the s coupling to the nucleon can be accidentally small, with a mild tuning of $O(10\%)$. The two cases described here correspond to the two points in Table I. The Sommerfeld enhancement factors for these points are $\gtrsim 100$ [8].

For the consistency of the DM story, $\tilde{s}$ must not be stable. In low-scale SUSY breaking, it is natural to assume that $\tilde{s}$ decays into the gravitino $\tilde{G}$. The mass of $\tilde{s}$ is typically above m_a, in which case the lifetime is given by $\tau_{\tilde{s}\to a\tilde{G}} \simeq 96\pi m_{3/2}^2 M_{\rm Pl}^2 / m_{\tilde{s}}^5$. For a gravitino mass $m_{3/2} \lesssim O(10-100~\text{eV})$, this is sufficiently short that $\tilde{s}$ never dominates the universe. Also, gravitinos this light do not cause a cosmological problem [29].

The light s, a, and $\tilde{s}$ states have interesting implications for LHC phenomenology. For example, the Higgs can decay as $h \to aa \to 4\mu$. Strongly produced SUSY particles will typically cascade decay into the light Higgsino, which will subsequently decay into $\tilde{s}$, sometimes by emitting an a or s. This leads to pairs and quartets of collimated muons with small invariant mass.

Finally, we note that in general, DM can either be the fermion or scalar component of Ψ/Ψ^c, depending on which is lighter. Scalar DM works similarly to fermion DM, except the dominant annihilation modes are now ss and aa. Since the scalar/fermion mass splitting controls radiative corrections to m_S^2, the Ψ/Ψ^c states could be nearly degenerate, leading to coannihilation at freezeout.

Outlook. We have presented a DM scenario that naturally explains the PAMELA and ATIC/PPB-BETS data. DM is a TeV-scale particle annihilating into an axion a, and the halo annihilation rate is enhanced through the scalar s. In the simplest realization, a is associated with a PQ symmetry, and we have constructed a corresponding SUSY model where the Higgsino and DM masses have a common origin from $U(1)_{\rm PQ}$ breaking.

There are other implementations of our scenario. For example, in models of low-scale dynamical SUSY breaking, the sector breaking SUSY typically leads to an R axion with the decay constant $f_a \sim \Lambda/4\pi$, where $\Lambda \approx O(10-100~\text{TeV})$ is the dynamical scale. This axion can serve as our a if the Higgses and DM obtain $U(1)_R$ breaking masses. The mass of a is $m_a^2 \sim \Lambda^3/M_{\rm Pl}$ [30], and the scales could work with $O(1)$ (or loop) factors.

One could also consider a purely leptonic axion, for example, by introducing a separate $U(1)_X$ and Higgs fields for the lepton sector. In this case a does not have hadronic couplings, eliminating the tension with direct detection experiments and opening the possibility for $a \to e^+e^-$ decays.

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[31] DM may be a Majorana fermion ψ obtaining a mass from $\mathcal{L} = -\xi S\psi^2/2+$h.c., in which case the stability is ensured by a Z_2 symmetry.
[32] Alternatively, we can make ψ/ψ^c charged under a new $U(1)_{\rm NE}$ gauge symmetry. As long as the $U(1)_{\rm NE}$ gauge coupling is small, $g_{\rm NE} \lesssim \xi$, thermal freezeout is unaffected and the halo boost is $B \sim \alpha_{\rm NE}/v_{\rm halo}$. However, $U(1)_{\rm NE}$ must be spontaneously broken to avoid constraints from big-bang nucleosynthesis, since axion decays would otherwise generate a large energy density of $U(1)_{\rm NE}$ gauge bosons. Nonperturbative enhancements through massless gauge bosons are also ruled out by astrophysical observations [18]. The model presented in the text avoids both bounds automatically, since the scalar s has a mass $\gtrsim O(1 \text{ GeV})$.
[33] The bounds from photon flux may be weakened by assuming a different DM density or velocity profile, or due to the uncertainty in Eq. (9). In this case, the region $2m_\tau < m_a < 2m_b$ may open up, since hadronic decays are suppressed by $3m_c^2/m_\tau^2 \tan^4\beta$. The bounds from Υ decay is satisfied; the predicted rate is $\text{Br}(\Upsilon \to \gamma a) \simeq 1 \times 10^{-4} \sin^4\beta\, (v_{\rm EW}/f_a)^2$, while the experimental bound is $\text{Br}(\Upsilon \to \gamma a) \sim 10^{-4}$ with $a \to \tau^+\tau^-$ [25]. The region $800 \text{ MeV} \lesssim m_a \lesssim 2 \text{ GeV}$ may also be allowed.