Superheavy dark matter from thermal inflation

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It is quite plausible that the mass of the dark matter particle increases significantly after its freeze-out, due to a scalar field rolling to large values. We describe a realization of this scenario in the context of thermal inflation which naturally gives a cold dark matter particle with the correct cosmological abundance and a mass around 10^{10} GeV, evading the conventional upper bound of 10^5 GeV. We also discuss another realization which could produce a cosmologically interesting abundance of near Planck mass, possibly electromagnetically charged, particles. The detection and observational consequences of superheavy cold dark matter or wimpzillas are briefly examined. [S0556-2821(99)04314-3]

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I. INTRODUCTION

A ''model-independent'' bound of about 10^5 GeV [1] has often been invoked (see [2] and references therein) as the largest possible mass the dark matter particle can have. The derivation of this bound uses the unitarity bound on the annihilation cross section, and makes one crucial assumption: thermal equilibrium in the early universe. The unitarity bound on the annihilation cross section tells us

$$\langle \sigma_A v \rangle \lesssim \frac{1}{M^2} \tag{1}$$

where M is the mass of the dark matter particle, and $\langle \sigma_A v \rangle$ is the thermally averaged annihilation cross section that appears in the relevant Boltzman equation [3] (i.e., annihilation per unit time is given by $n\langle \sigma_A v \rangle$, where n is the proper number density of the dark matter particle). The assumption of thermal equilibrium, on the other hand, tells us that the freeze-out abundance is given by the thermal distribution

$$n \sim (MT_f)^{3/2} \exp\left(-\frac{M}{T_f}\right) \tag{2}$$

where T_f is the freeze-out temperature.[1] The exponential suppression implies that T_f cannot be too much smaller than M; hence, combining this with the unitarity bound, it can be shown that [1]

$$\Omega h^2 \gtrsim 0.1 \left(\frac{M}{10^5 \ \mathrm{GeV}}\right)^2 \tag{3}$$

where Ω is the ratio of the mass density in the dark matter particle to the critical density today, and h is the Hubble constant in units of 100 km $\mathrm{s}^{-1}\mathrm{Mpc}^{-1}$. This bound naturally has important implications for dark matter searches.

Recently, it has been shown that this bound can be evaded by violating the assumption of thermal equilibrium in the early universe, for instance by producing the dark matter particle at the end of inflation via preheating or reheating or gravitational particle creation [4,5]. The masses required for a present abundance of $\Omega \sim 1$ are within a few orders of magnitude of the Hubble parameter at the end of inflation, i.e. $\sim 10^{12} - 10^{16}$ GeV. For this reason, they have been called superheavy dark matter or wimpzillas.

Here, we propose a different production mechanism that can also evade the 10^5 GeV upper bound, and naturally achieves $\Omega \sim 1$. It makes use of a late period of inflation called thermal inflation [6–8] that occurs at an energy scale of around 10^6 GeV, with a Hubble parameter of the order of 1 keV, which is to be compared with the GUT-scale ordinary inflation having $V^{1/4} \sim 10^{16}$ GeV and $H \sim 10^{13}$ GeV used in [4,5].

Thermal inflation provides a natural solution to the Polonyi or moduli problem [9] that generically arises in string theory (see also [10]). It occurs when a ''flaton,'' a scalar field with a small mass and a large vacuum expectation value, is held at the origin by its finite temperature effective potential. One gets a few e-folds of inflation because the flaton's potential energy dominates the thermal energy density well before the temperature drops below the critical temperature for the flaton to start rolling away from the origin. The prototypical flaton potential is described in Sec. II A. Thermal inflation occurs at a very low energy scale which is the reason it can successfully dilute the potentially harmful moduli produced after ordinary inflation. Note that it will also dilute any superheavy dark matter produced at the end of ordinary inflation.

Our idea for dark matter production works roughly as follows. A particle ψ and its antiparticle $\bar{\psi}$, which carry some conserved charge to make them stable, are coupled to the flaton. They are initially massless during the thermal inflation when the flaton is held at the origin by the finite temperature effects of ψ and $\bar{\psi}$. After the temperature of the universe drops below the critical temperature, the flaton begins to fast-roll down its potential. The ψ particle quickly gains mass in the process, which reduces its annihilation cross section, and its abundance quickly freezes out. After

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[1]This assumes a cold relic. For a hot relic that freezes out at $T_f \gtrsim M$, the number density will be higher leading to a stronger bound on its mass [3].

that, the flaton continues to roll until it reaches the true vacuum, acquiring a large expectation value and giving ψ a large mass.[2] The parameters for thermal inflation give a mass for ψ of around 10^{10} GeV, and work out naturally to give an abundance of $\Omega_{\psi\bar{\psi}}\sim 1$. The details are explained in Sec. II.

This production mechanism evades the conventional upper bound of 10^5 GeV by giving the particle ψ a larger annihilation cross section at freeze-out than what one would expect based on its final mass, and by entropy production after the freeze-out.

Since thermal inflation provides a natural solution to the moduli problem, and since the prediction of $\Omega_{\psi\bar{\psi}}\sim 1$ follows rather naturally from its parameters (assuming ψ is stable), the possibility of a significant fraction of the universe being composed of superheavy dark matter should be taken seriously. In Sec. III, we discuss the observational consequences and detectability of such dark matter, which could be completely inert, weakly-interacting, or even electromagnetically or strongly-charged. Finally, we conclude in Sec. IV with discussions of other plausible realizations of the mechanism outlined above, which include the possibility of producing near Planck mass relics that could perhaps be electromagnetically charged.

II. SUPERHEAVY DARK MATTER FROM THERMAL INFLATION

A. Particle physics

A simple model of thermal inflation is provided by the superpotential [12]

$$W(\phi)=\frac{\lambda_\phi\phi^4}{4M_{\rm Pl}} \tag{4}$$

where ϕ is the flaton, and $M_{\rm Pl}\simeq 2.4\times 10^{18}$ GeV. This form for the superpotential can be guaranteed by a Z_4 discrete gauge symmetry because the superpotential is a holomorphic function of ϕ, i.e., does not depend on ϕ's complex conjugate $\phi^\dagger$. The supersymmetric part of the scalar potential is then given by

$$V_{\rm susy}(\phi,\phi^\dagger)=\left|\frac{\partial W}{\partial\phi}\right|^2=\frac{|\lambda_\phi|^2|\phi|^6}{M_{\rm Pl}^2}. \tag{5}$$

In addition, one requires ϕ's soft supersymmetry-breaking mass-squared to be negative. The scalar potential is then

$$V(\phi,\phi^\dagger)=V_{\rm ti}-m_\phi^2|\phi|^2+\left(\frac{A\lambda_\phi\phi^4}{M_{\rm Pl}}+{\rm c.c.}\right)+\frac{|\lambda_\phi|^2|\phi|^6}{M_{\rm Pl}^2} \tag{6}$$

where $V_{\rm ti}$ and A are other soft supersymmetry-breaking parameters. One would expect $|A|\lesssim m_\phi$. The scale of the soft supersymmetry-breaking parameters of the minimal supersymmetric standard model (MSSM) is expected to be around 100 GeV to 1 TeV. This potential has four degenerate minima[3] with $|\phi|=\phi_{\rm vac}$, where

$$\phi_{\rm vac}^2=\frac{m_\phi M_{\rm Pl}}{\sqrt{3}\,|\lambda_\phi|}\left(\sqrt{1+\frac{4|A|^2}{3m_\phi^2}}+\frac{2|A|}{\sqrt{3}\,m_\phi}\right). \tag{7}$$

For $m_\phi=100$ GeV, $|\lambda_\phi|=1$, and neglecting A, this gives $\phi_{\rm vac}=10^{10}$ GeV. Requiring zero cosmological constant at the minima gives

$$V_{\rm ti}=\frac{2}{3}m_\phi^2\phi_{\rm vac}^2\left[1+\frac{|A|}{\sqrt{3}\,m_\phi}\left(\sqrt{1+\frac{4|A|^2}{3m_\phi^2}}+\frac{2|A|}{\sqrt{3}\,m_\phi}\right)\right]. \tag{8}$$

Thermal inflation starts when the energy density of the universe starts to be dominated by $V_{\rm ti}$, and ends a few e-folds later when the temperature drops below that required to hold ϕ at $\phi=0$. ϕ then rapidly rolls towards, and oscillates about, the minima of its potential. It eventually decays, leaving a radiation dominated universe, at a temperature $T_{\rm dec}$. We require

$$T_{\rm dec}\gtrsim 10\ {\rm MeV} \tag{9}$$

to avoid interfering with nucleosynthesis.

In order for ϕ to be held at $\phi=0$ by thermal effects during the thermal inflation, ϕ must have unsuppressed interactions with other fields in the thermal bath. We will assume these interactions include a coupling of the form[4]

$$W=\lambda_\psi\phi\bar{\psi}\psi \tag{10}$$

with $|\lambda_\psi|\sim 1$. After thermal inflation, ψ and $\bar{\psi}$ will acquire masses

$$M=|\lambda_\psi|\,\phi_{\rm vac} \tag{11}$$

from this coupling. For $|\lambda_\psi|=1$ and $\phi_{\rm vac}=10^{10}$ GeV, this gives $M=10^{10}$ GeV.

In order to satisfy the decay constraint, Eq. (9), we require ϕ to have some, possibly indirect, couplings to the MSSM, beyond the ever present gravitational couplings. We will assume this is achieved by having ψ and $\bar{\psi}$ charged under ${\rm SU(3)}\times{\rm SU(2)}\times{\rm U(1)}$. Other possibilities were considered in Ref. [7]. In order not to interfere with gauge coupling

[2] A related mechanism has also been considered in [11].

[3] The domain walls associated with this degeneracy are most likely harmless, either because the vacua are identified because of a discrete gauge symmetry, or because higher order terms, that would generically be present in the absence of a gauge symmetry, break the degeneracy.

[4] A flaton ϕ cannot have a coupling of the form $W\sim\phi^2\psi$ since it would lead to an unsuppressed quartic self-coupling in the potential for ϕ. Such a coupling can be forbidden by an appropriate gauge symmetry. The only other possibility would be for ϕ to be charged under some continuous gauge symmetry, which in the vacuum would be broken at a scale $\sim\phi_{\rm vac}\sim 10^{10}$ GeV. We do not consider this possibility here.

unification, ψ and $\bar{\psi}$ should form complete representations of SU(5). These couplings give a decay rate [13]

$$\Gamma_\phi \sim 3\times 10^{-9}\left(\frac{g_\psi^2 m_\phi^3}{\phi_{\rm vac}^2}\right) \tag{12}$$

where g_ψ is the number of internal degrees of freedom of ψ and $\bar{\psi}$. For example, if $\psi=\mathbf{5}$ and $\bar{\psi}=\overline{\mathbf{5}}$ then $g_\psi=40$. The decay temperature is therefore [3]

$$T_{\rm dec}\simeq g_*^{-1/4}\Gamma_\phi^{1/2}M_{\rm Pl}^{1/2}\sim 2\times 10^{-5}\left(\frac{g_\psi m_\phi^{3/2}M_{\rm Pl}^{1/2}}{\phi_{\rm vac}}\right)$$
$$\sim 300\,{\rm MeV}\left(\frac{g_\psi}{100}\right)\left(\frac{m_\phi}{100\,{\rm GeV}}\right)^{3/2}\left(\frac{10^{10}\,{\rm GeV}}{\phi_{\rm vac}}\right) \tag{13}$$

where g_* is the effective number of massless degrees of freedom in the universe at temperature $T_{\rm dec}$. Note that parametric resonance is unlikely to be important because ϕ oscillates around a large vacuum expectation value rather than the origin.

Renormalization, amongst other things, will split the degeneracy of Eq. (11) amongst the various components of ψ and $\bar{\psi}$. The renormalization will tend to make SU(3) charged components the heaviest, and SU(3)$\times$SU(2)$\times$U(1) singlet components, if they exist, the lightest [14].

Our main unjustified assumption is now to assume that ψ and $\bar{\psi}$ carry opposite charge under some discrete (or continuous) gauge symmetry under which all other fields are neutral. This symmetry is, however, helpful in avoiding unwanted superpotential couplings to MSSM fields. The lightest component of ψ and $\bar{\psi}$ will then be absolutely stable and so potentially a dark matter candidate.[5] For example, we could have a Z_8 discrete gauge symmetry, under which ϕ, ψ, and $\bar{\psi}$ have charges 2, -1, and -1, respectively. This would guarantee Eqs. (4) and (10), and after ϕ acquires its vacuum expectation value, the Z_8 will be broken down to a Z_2 under which only ψ and $\bar{\psi}$ are charged. A simple choice of representations that satisfies the discrete anomaly cancellation conditions [15] is $\psi=\mathbf{16}+\mathbf{1}$ and $\bar{\psi}=\overline{\mathbf{16}}+\mathbf{1}$ of SO(10).

B. Abundance

During thermal inflation, ψ and $\bar{\psi}$ will be relativistic and in thermal equilibrium. Their number density will therefore be

$$n(T)=\frac{7\zeta(3)g_\psi T^3}{8\pi^2} \tag{14}$$

where $\zeta(3)\simeq 1.202$, and g_ψ is the number of internal degrees of freedom of ψ and $\bar{\psi}$.[6] If $\psi=\mathbf{16}+\mathbf{1}$ and $\bar{\psi}=\overline{\mathbf{16}}+\mathbf{1}$ of SO(10), as in the example of the previous section, then $g_\psi=136$. Through the coupling of Eq. (10), ψ and $\bar{\psi}$ will generate the finite temperature effective potential for ϕ

$$V_T(\phi)=V_{\rm ti}+\left(\frac{g_\psi}{32}\left|\lambda_\psi\right|^2T^2-m_\phi^2\right)|\phi|^2+\cdots. \tag{15}$$

Thermal inflation will therefore end at the temperature

$$T_{\rm c}=\frac{4\sqrt{2}\,m_\phi}{\sqrt{g_\psi}\,|\lambda_\psi|} \tag{16}$$

when ϕ begins to roll away from $\phi=0$. Shortly afterwards, the abundance of ψ and $\bar{\psi}$ freezes out. Meanwhile, ϕ will continue to roll towards its vacuum expectation value $|\phi|=\phi_{\rm vac}$. Once ϕ acquires its vacuum expectation value $|\phi|=\phi_{\rm vac}$, the coupling of Eq. (10) will give ψ and $\bar{\psi}$ masses

$$M=|\lambda_\psi|\,\phi_{\rm vac}. \tag{17}$$

The freeze-out abundance of ψ and $\bar{\psi}$ can be estimated as follows. First, it is important to keep in mind that the temperature does not drop as ϕ rolls from 0 to $\phi_{\rm vac}$. This is because the time-scale for the roll-over is m_ϕ^{-1}, which is much smaller than the Hubble time $H_{\rm ti}^{-1}\sim m_\phi^{-1}\phi_{\rm vac}^{-1}M_{\rm pl}\gg m_\phi^{-1}$. Hence $T=T_{\rm c}$ throughout the roll-over. The freeze-out occurs as ψ gains mass and begins to become non-relativistic when its thermal abundance is given by

$$n(m_\psi)=g_\psi\left(\frac{m_\psi T_{\rm c}}{2\pi}\right)^{3/2}\exp\left(-\frac{m_\psi}{T_{\rm c}}\right) \tag{18}$$

where m_ψ denotes the ϕ dependent mass of ψ, $m_\psi=|\lambda_\psi\phi|$, which increases as ϕ rolls down the potential. In other words, in contrast with the usual freeze-out calculation, it is m_ψ that is changing with time rather than the temperature. The freeze-out abundance is determined by equating the annihilation rate $\Gamma_{\psi\bar{\psi}}$ with the inverse-time-scale of the problem at hand, i.e., m_ϕ

$$\Gamma_{\psi\bar{\psi}}(m_\psi)\sim n(m_\psi)\langle\sigma|v|\rangle(m_\psi)\sim m_\phi. \tag{19}$$

Using the fact that $\langle\sigma|v|\rangle\lesssim 1/m_\psi^2$, and that $m_\phi\sim|\lambda_\psi|T_{\rm c}$, it is not hard to see that the freeze-out occurs when $m_\psi\sim T_{\rm c}$, and that the freeze-out abundance is given by Eq. (14) with $T=T_{\rm c}$, suppressed by at most a factor $|\lambda_\psi|$. So for $|\lambda_\psi|\sim 1$, and to within an order of magnitude, there is no significant net annihilation of ψ and $\bar{\psi}$ after the beginning of the roll-over, and the freeze-out occurs well before ϕ reaches $\phi_{\rm vac}$.

After the freeze-out, ϕ will continue to roll towards its vacuum expectation value, increasing the mass of ψ to a

[5]Throughout this paper, we use the same notation ψ for the complete representation and its lightest component (the dark matter).

[6]We have assumed, as is appropriate for a supersymmetric theory, that there are equal numbers of bosonic and fermionic degrees of freedom [7/8=(1+3/4)/2].

large value. It can be checked that the annihilation rate by the time ϕ reaches the minimum is negligible[7]

$$\Gamma_{\psi\bar{\psi}} \lesssim \frac{T_c^3}{M^2} \sim \frac{m_\phi^3}{|\lambda_\psi|^5 \phi_{\rm vac}^2} \sim \left(\frac{|\lambda_\phi|^{3/2} m_\phi^{1/2}}{|\lambda_\psi|^5 M_{\rm Pl}^{1/2}}\right) H_{\rm ti} \ll H_{\rm ti}. \quad (20)$$

Subsequently, ϕ will oscillate around $\phi_{\rm vac}$.[8]

One might worry that during the oscillation, a significant amount of annihilation or production of ψ and $\bar{\psi}$ might occur if ϕ returns to small values. The rapid buildup of gradient energy will prevent ϕ from returning to small values except in a few isolated places. In addition, ϕ would eventually be prevented from returning to small values by the Hubble expansion. Strings with walls attached are also formed, which will likely disappear quickly [16]. Their direct radiation into the heavy ψ particles is heavily suppressed because $M_\psi \gg m_\phi$. On the other hand, the ψ particles are light in the cores of the strings, and could be created and trapped there. If a string loop carries a net ψ charge, it will be released in the form of (heavy) ψ particles when the loop annihilates. If we assume each string produces of the order of one ψ particle, our rough estimate of the ψ abundance should still be valid.

The energy density in ψ and $\bar{\psi}$ will then scale with that of the oscillating flaton until the flaton finally decays, leaving a radiation dominated universe at a temperature $T_{\rm dec}$. The energy density of ψ and $\bar{\psi}$ will then scale with the entropy density of the universe, s. The current value of the entropy density is

$$s_0 = 2.2 \times 10^{-38}\,{\rm GeV}^3. \quad (21)$$

Finally, we wish to compare the current energy density of ψ and $\bar{\psi}$ with the critical density

$$3H_0^2 = 3 \times 10^{-47}\,{\rm GeV}^4. \quad (22)$$

For ψ and $\bar{\psi}$ to be a viable dark matter candidate we require

$$\Omega_{\psi\bar{\psi}} \equiv \frac{\rho_{\psi\bar{\psi}}}{3H_0^2} \sim 0.3. \quad (23)$$

Putting everything together we get

$$\Omega_{\psi\bar{\psi}} = n(T_c) M \left(\frac{\rho_{\rm dec}}{V_{\rm ti}}\right)\left(\frac{s_0}{s_{\rm dec}}\right)\left(\frac{1}{3H_0^2}\right) \quad (24)$$

where $\rho_{\rm dec}$ is the energy density of the universe at the end of the flaton decay. Therefore, using $\rho_{\rm dec} = \frac{3}{4} T_{\rm dec} s_{\rm dec}$,

$$\Omega_{\psi\bar{\psi}} = |\lambda_\phi|\,|\lambda_\psi|^{-2} \left(\frac{g_\psi}{100}\right)^{1/2} \left(\frac{m_\phi}{100\,{\rm GeV}}\right)^{3/2} \left(\frac{5 \times 10^4 \phi_{\rm vac} T_{\rm dec}}{g_\psi m_\phi^{3/2} M_{\rm Pl}^{1/2}}\right) \times \left(\frac{8\pi^2 n(T_c)}{7\zeta(3) g_\psi T_c^3}\right) \left(\frac{\sqrt{g_\psi}\,|\lambda_\psi| T_c}{4\sqrt{2}\, m_\phi}\right)^3 \left(\frac{M}{|\lambda_\psi|\,\phi_{\rm vac}}\right) \times \left(\frac{2 m_\phi^2 \phi_{\rm vac}^2}{3 V_{\rm ti}}\right) \left(\frac{m_\phi M_{\rm Pl}}{\sqrt{3}\,|\lambda_\phi|\,\phi_{\rm vac}^2}\right) \quad (25)$$

where all factors in brackets are of order 1. We have displayed explicitly the assumed relations between the various quantities, e.g., $M = |\lambda_\psi|\,\phi_{\rm vac}$, etc.

III. OBSERVATIONAL CONSTRAINTS AND DETECTION

The above simple model leaves open the question of what kind of interaction ψ has with ordinary matter. The same is true of other production mechanisms of superheavy dark matter [4,5]. Let us go through the different possibilities one by one.

Electromagnetically charged. These have been referred to as charged massive particles (CHAMPs) in the literature [17]. At late times, they primarily take the form of $p^+\psi^-$ (hydrogen with a heavy ''electron'' which has very low cross section with other atoms), $\psi^+ e^-$ (heavy hydrogen) or $\psi^- {\rm He}^{++} e^-$ (ψ^- bound to the helium nucleus to make another kind of heavy hydrogen). Various constraints exist on such particles, ranging from the absence of heavy-hydrogen-like atoms in water to nondetection in γ-ray and cosmic-ray detectors [17]. By far, the strongest constraint appears to come from the existence of old neutron stars [18], where only $M \gtrsim 10^{16}\,{\rm GeV}$ is allowed. Otherwise, a sufficient net number of ψ^+ particles collects in the neutron star, forms a black hole in the center and eats up the star on a short time scale; this is in part because the hydrogen-heavy-hydrogen scattering cross section is high, given by the square of the Bohr radius in the low velocity limit $\sigma \sim 10^{-17}\,{\rm cm}^{-2}$. However, if the abundance of ψ^+ and that of ψ^- bound to helium are the same in the halo, the constraint is weakened to $M \gtrsim 10^{10}\,{\rm GeV}$. The conventional bound of $M \lesssim 10^5\,{\rm GeV}$ would then be fatal to the existence of such particles. The kind of production mechanism like the one proposed here, or elsewhere, which evades the conventional bound, could resurrect the intriguing idea that the dark matter can be charged. But as we can see, significant astrophysical constraints already exist. It should be noted that the near Planck mass relic that will be discussed in Sec. IV satisfies even the demanding bound of $M \gtrsim 10^{16}\,{\rm GeV}$. The economic importance of such stable massive electromagnetically charged particles cannot be overestimated. For example, they could be used to catalyze nuclear fusion [17].

Strongly charged. These have been referred to as strongly interacting massive particles (SIMPs) in the literature [19,20]. Significant bounds on their masses, if they have significant cosmological abundances, come from nucleosynthe-

[7] If $|\lambda_\psi|$ is small, the fifth power of $|\lambda_\psi|$ in this formula could make $\Gamma_{\psi\bar{\psi}}$ significant which means that there would be some annihilation of ψ and $\bar{\psi}$. However, in this case the cosmological abundance will also be boosted up by a higher T_c [see Eqs. (16) and (25)].

[8] Note that the backreaction of the finite density of ψ particles on ϕ's potential is negligible.

sis as well as the absence of anomalously heavy isotopes of familiar nuclei [21]. A systematic study of constraints from direct detection and the existence of old neutron stars and the Earth was made in Ref. [19]. Assuming there is no ψ-$\bar{\psi}$ asymmetry, the neutron-star argument and underground plus balloon experiments provide competitive bounds: only $M \gtrsim 10^8 - 10^{10}\,\mathrm{GeV}$ is allowed for a ψ-proton cross section of $\sigma \sim 10^{-30} - 10^{-25}\,\mathrm{cm}^{-2}$. As in the case of electromagnetically charged dark matter, the possibility of producing them without overclosing the universe gives such dark matter candidates a new life.

Weakly charged. Naturally, no significant constraints exist if ψ has only weak-scale interactions like the neutralino [i.e., $\sigma \sim 10^{-44}(m_n/\mathrm{GeV})^4\,\mathrm{cm}^2$ in the large M limit, where m_n is the mass of the relevant nucleon]. Direct detection appears rather difficult simply because the halo number density scales as M^{-1}, and the neutralino with mass $\sim 100\,\mathrm{GeV}$ is already difficult to detect. Note that a large mass does not increase significantly the nuclear recoil: $\Delta E \sim M^2 m_n v^2/(M+m_n)^2$ where $v \sim 200\,\mathrm{km\ s}^{-1}$ is the average halo velocity of these particles; in the large M limit, ΔE is asymptotically M-independent. Indirect detection might seem even more hopeless. Not only does the halo number density drop by a factor of M, the annihilation (which gives rise to neutrinos) rate is suppressed by M^2 according to the unitary bound. However, three opposing factors help us here. First, on a sufficiently long time scale, the neutrino flux from ψ-$\bar{\psi}$ annihilation in the core of the Sun or the Earth is determined not by the annihilation rate, but by the capture rate, which depends on the scattering cross section with nucleons and is not heavily suppressed. Second, indirect detection works by observing muons that result from the interaction of the neutrinos with the rocks of the Earth. The cross section for producing muons and the range of the muons both scale up with energy, and hence the mass of ψ. A detailed calculation taking into account these effects as well as other relevant ones will be presented in a forthcoming paper.

Completely neutral. ψ could be a singlet under $\mathrm{SU}(3)\times\mathrm{SU}(2)\times\mathrm{U}(1)$. It is of course virtually impossible to detect such particles, except by their gravitational effects.

Lastly, superheavy dark matter has been postulated as the origin of ultra-high-energy cosmic-ray events at energies $\gtrsim 10^{10}\,\mathrm{GeV}$, above the Greisen-Zatsepin-Kuzmin cutoff [4,5]. The parameters of the model presented in the last section are sufficiently flexible to allow a mass of $M \sim 10^{11}\,\mathrm{GeV}$ to explain such events. The ψ particles cannot by themselves be the primaries because of the large mass, even if they have hadronic interactions [22,21]. The simplest way is to have them decay into hadrons, e.g., protons, which reach the Earth's atmosphere. But then, one has to invoke special reasons to explain why they are not stable but sufficiently long-lived.

IV. DISCUSSION

We have shown that a simple and well-motivated model of thermal inflation naturally produces dark matter particles ψ and $\bar{\psi}$ of mass $M \sim 10^{10}\,\mathrm{GeV}$ with a cosmological abundance in the correct range. As our mechanism is implemented by thermal inflation, which solves the moduli problem by late entropy production, it is robust against such dilution.

The same general mechanism could also be applied to what one might call ''moduli thermal inflation'' to produce near Planck mass particles at low energy scales. Moduli thermal inflation is a limit of thermal inflation in which the flaton is replaced by a modulus, so that roughly speaking $|\lambda_\phi| \sim m_\phi/M_{\mathrm{Pl}}$ in Eq. (6). ϕ_{vac} will then be of the order of the Planck scale. In the context of string theory, it is then very reasonable to assume that the vacuum expectation value of the modulus, $\phi = \phi_{\mathrm{vac}}$, corresponds to another ''origin'' in field space where new fields become light; for example one could have superpotential couplings of the form $W = \lambda_\chi(\phi - \phi_{\mathrm{vac}})\bar{\chi}\chi$. Such a ''coupled'' modulus would appear like an ordinary scalar field (e.g., a squark or slepton field) in the true vacuum.[9] The decay temperature would then no longer scale as in Eq. (13) but instead could be as high as $V_{\mathrm{ti}}^{1/4}$. Putting these modifications into Eq. (25) would also give us a value of $\Omega_{\bar{\psi}\psi}$ in the neighborhood of 1. However, another consequence of having $\phi_{\mathrm{vac}} \sim M_{\mathrm{Pl}}$ is that one could get a significant number of e-folds of non-slow-roll inflation as ϕ rolls from $\phi \sim 0$ to $\phi \sim \phi_{\mathrm{vac}}$. To avoid this inflation diluting the ψ particles too much, one would require $m_\phi \gtrsim 10 V_{\mathrm{ti}}^{1/2}/M_{\mathrm{Pl}}$ and so $\phi_{\mathrm{vac}} \lesssim M_{\mathrm{Pl}}/10$. This would limit the final mass of the ψ particles to be $M \lesssim \mathrm{few} \times 10^{17}\,\mathrm{GeV}$. However, this is still above even the stringent limit on electromagnetically charged dark matter obtained in Ref. [18].

Note that because moduli thermal inflation occurs at too high an energy scale to solve the moduli problem, this scenario would only be viable if there were no moduli problem[10] because otherwise the ψ particles would be diluted by another epoch of thermal inflation, or some other late entropy production, that would be needed to dilute the decoupled moduli produced at the end of the moduli thermal inflation.

A related scenario could emerge from some of the more plausible models of inflation [24]. Here ψ or $\bar{\psi}$ is the inflaton, which holds ϕ at zero by the hybrid inflation mechanism [25] rather than thermal effects. One could then even get dark matter in the form of charged, near Planck mass, inflatons.

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[9] It would be an excellent candidate for an Affleck-Dine field [23].

[10] For example, because all the moduli are of this coupled type, rather than the decoupled type that give rise to the moduli problem. How one fits the dilaton into such a picture is unclear though.

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