Opening the window on strongly interacting dark matter

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We discuss the possibility that the dark matter consists of strongly interacting massive particles (SIMP's) which have cross sections with ordinary matter which are larger than characteristic weak-interaction cross sections. We show that, while results from $\beta\beta$ decay, cosmic-ray detectors, galactic-halo stability, the cooling of molecular clouds, proton-decay detectors, and the existence of old neutron stars and the Earth constrain the interactions of the missing matter with ordinary matter over a broad range of parameter space, there still exist several windows for SIMP's. It is noteworthy that there are two regions of less than geometric cross sections: one with masses of 10^5-10^7 GeV and another with masses above 10^{10} GeV.

I. INTRODUCTION

Since the idea of missing matter was introduced, it has always been assumed that its interactions with ordinary matter are weak—at least as small as ordinary weak-interaction cross sections. This prejudice has received scant critical evaluation. There has been only one experiment[1] which focused on probing the regime of strongly interacting massive particles (SIMP's). However, the recent suggestion by Dimopoulos, Eichler, Esmailzadeh, and Starkman[2] and by De Rújula, Glashow, and Sarid[3] that the dark matter could be charged has led to an examination by these groups as well as others of the existing constraints, experimental and astrophysical, on the cross sections of the "dark matter" on ordinary matter. In this paper we present new limits on the dark-ordinary transfer cross section $\sigma_{\rm tr}$ as a function of the mass m_h of the dark-matter particle. These limits are mostly obtained by modifying the arguments originally applied to charged massive particles. Specifically, they come from cosmic-ray plastic track and solid-state detectors, $\beta\beta$-decay spectrometers, infall of the galactic halo, heating of molecular clouds, the longevity of neutron stars and of the Earth, proton-decay detectors, and heavy-isotope searches. We show that, although broad regions of parameter space ($\sigma_{\rm tr}$ vs m_h) are excluded, there exist four large allowed regions for SIMP's. In two of these regions (10^5 GeV $\lesssim m_h \lesssim 10^7$ GeV and $m_h \gtrsim 10^{10}$ GeV) the cross sections on nuclei are less than geometric, while in the other two (3 GeV $\lesssim m_h \lesssim 10^4$ GeV and $m_h \gtrsim 10^{10}$ GeV) they are greater than geometric. (The last region is allowed only for so-called "coherent" interactions.) In addition, the regions have complicated structures which allow only either fully (particle-antiparticle) symmetric or fully asymmetric SIMP abundances in certain subregions.

Throughout this paper we consider only the galactic-halo dark matter. We take the local mass density of the halo to be $\rho=0.4$ (GeV/c^2) cm^{-3}, and assume that the dark matter is not clumped.

II. LIMITS ON $\sigma_{\rm tr}$

We now proceed to discuss the various limits. In Fig. 1 the raw limits are plotted, with no attempt to relate the cross sections on different materials. In Figs. 2 and 3 we plot the limits on the dark-matter (transfer) cross section on protons σ_p, assuming that $\sigma(A,Z)=A^2[m_{\rm red}(A)/m_{\rm red}(1)]^2\sigma_p$ and $\sigma(A,Z)=\text{spin}(A,Z)[m_{\rm red}(A)/m_{\rm red}(1)]^2\sigma_p$, corresponding to coherent (spin-independent) and spin-dependent cross sections. Here spin (A,Z) is one if the nucleus has spin, and zero otherwise. We show only the region $m \geq 1$ GeV, since below this value essentially all of our limits disappear. We also do not graph the allowed region which obviously exists at small σ, as this is not the subject of this paper.

A. Semiconductor dark-matter detectors

Detectors for halo dark matter work on the principle of detecting the energy deposited by nuclei recoiling from elastic collisions with halo particles. Semiconductor devices, such as the Ge $\beta\beta$-decay spectrometers[4,5] or silicon detectors[6] convert the nuclear recoils to electron-hole pairs, which can be amplified and detected. Since the $\beta\beta$-decay spectrometers were designed for ultralow background experiments, they require shielding from cosmic rays and have been placed at several hundred to thousand meters of water equivalent (m.w.e.) below the Earth's surface. This unfortunately reduces their utility for searching for dark matter with a large $\sigma_{\rm tr}$, since the overlying

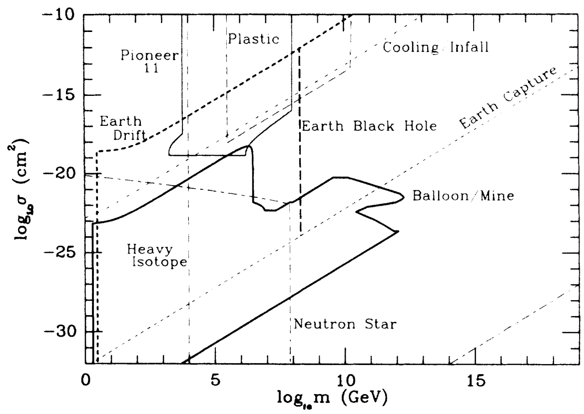

FIG. 1. The limits on dark-matter scattering cross sections as function of mass coming from balloon-borne and underground experiments (bold solid), galactic-halo infall and cloud cooling (dotted-dashed), Pioneer 11 (dotted), plastic track detectors (light long-dashed), the existence of old neutron stars (long-short-dashed), the absence of a neutrino signal for particles which drift into Earth's core (bold short-dashed and light short-dashed), the absence of a black hole in Earth (bold long-dashed), and the absence of a ultraheavy isotopes of light elements (dotted-dashed). Bold lines exclude down and to the right, light lines upward. The "Heavy Isotope" line excludes to the left. No assumptions are made relating the various cross sections to each other.

Earth degrades the SIMP energy. Nevertheless, their low backgrounds will prove useful in extending the limits to low flux, i.e., high m_h. A balloon-borne experiment[6] by Rich, Rocchia, and Spiro was designed to extend the excluded region to higher cross sections; however, since it necessarily has higher backgrounds it becomes flux limited at lower masses than do the underground detectors.

The energy of a SIMP as it passes through material is degraded according to

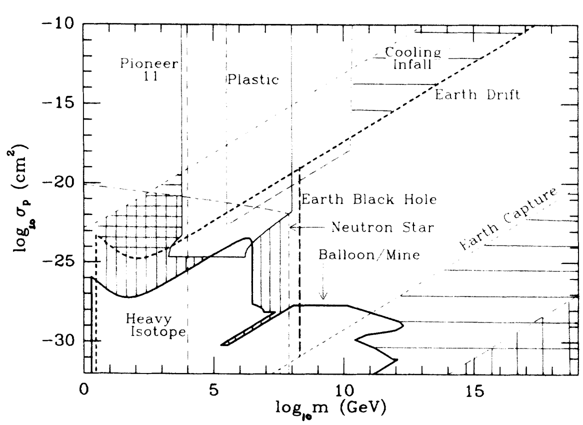

FIG. 2. The limits on the dark-matter cross section on protons as a function of mass, assuming that $\sigma(A,Z) = A^2[m_{\rm red}(A)/m_{\rm red}(1)]^2\sigma_p$. The coding of the lines is as in Fig. 1. The vertically hatched regions are allowed only with an asymmetry, the horizontally hatched regions are allowed only in the absence of an asymmetry, and the cross-hatched regions are unconstrained.

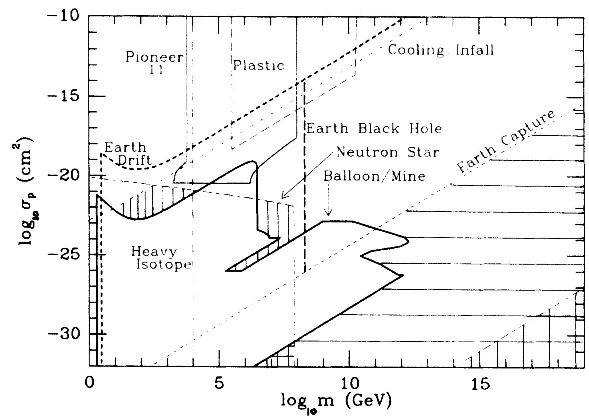

FIG. 3. The limits of the dark-matter cross section on protons as functions of mass, assuming that $\sigma(A,Z) = \text{spin}(A,Z)[m_{\rm red}(A)/m_{\rm red}(1)]^2\sigma_p$, where spin is one for odd A and zero otherwise. The coding of the lines and the shading of allowed regions are as in Fig. 2.

$$\frac{dE}{dx} = -\sum_A n_A \sigma_A \left[\frac{m_{\rm red}(A)^2}{m_A} \right] v^2 , \tag{2.1}$$

where n_A is the number density of a given isotope, m_A is the mass of the isotope, $m_{\rm red}(A)$ is the reduced mass of the isotope-SIMP system, and v is their relative velocity. $m_{\rm red}(A)^2 v^2/m_A$ is the average energy transfer per collision assuming isotropic scattering. Thus, at a mass column density of (ρx),

$$E(\rho x) = E(0)\exp\left[-\frac{2\sigma_p}{m_h}(\rho x)\mathcal{F} \right] , \tag{2.2}$$

where σ_p is the cross section on hydrogen and

$$\mathcal{F} = \sum_A f_A \left[\frac{m_{\rm red}(A)}{m_A} \right]^2 \frac{\sigma_A}{\sigma_p} .$$

Here f_A is the mass fraction of a given isotope.

At the upper limits of the cross sections of interest, a SIMP has a mean free path much smaller than the size of these detectors. Therefore, should it reach the detector with sufficient energy to cause a detectable total nuclear recoil event ($E_{\rm recoil} \gtrsim 2$ keV for Si, $\gtrsim 4$ keV for Ge), the particle will be detected. (Since the SIMP will actually interact many times in the detector, and hence deposit many times $E_{\rm recoil}$ overall, the criterion actually has to do with having a significant conversion efficiency for the nuclear recoil energy into electron-hole pairs. Moreover, since the SIMP is depositing so much energy, it is necessary that the detector be sensitive to events which contain MeV's or GeV's of energy. In the balloon detectors as designed, these events therefore register in the overflow bin.) Since $E_{\rm recoil} \sim [2m_{\rm red}(A)/m_h]E_h$ we can determine the minimum energy $E_{\rm min}$ to which the h can be degraded and still be detectable. Knowing the depth (ρx) of the detector and the composition of the overlying matter we will be able to set a limit of

$$\frac{\sigma_H}{m_h} > \frac{1}{2}\ln\left[\frac{E(0)}{E_{\min}}\right]\frac{1}{\rho x \mathcal{F}} \tag{2.3}$$

on the transfer cross section.

In the balloon-borne experiment, data were taken from 100 m above ground up to a float level of 4.5 g/cm^2 (to which we must add the contribution of a 0.5-mm quartz window having a column density of 1.1 g/cm^2). Allowing an energy degradation of 90%, so that $E_{\rm recoil} \gtrsim 2.7$ keV, we find $\sigma_H/m_h \gtrsim 3.6\times10^{-25}/\mathcal{F}$ cm^2/GeV for the high-altitude data, $\sigma_H/m_h \gtrsim 1.8\times10^{-27}/\mathcal{F}$ cm^2/GeV for the low-altitude data. We can interpret these as limits on the "raw" cross section by setting $A=\overline{A}$, the average atomic number of the atmosphere plus window, so that $\mathcal{F}=[m_{\rm red}(\overline{A})/m_{\overline{A}}]^2$; or we can compute $\mathcal{F}$ for coherent and spin-dependent cross sections appropriately weighted for the composition of the overlying material.

As we shall see below, due to the cosmic-ray background the balloon-borne experiment is only applicable up to approximately 10^6 GeV. To push beyond this mass, one must consider the underground experiments. That of Caldwell *et al.* is the most obviously suited to our purposes, since it is located below a dam, only 600 m.w.e. below ground. However, because these experiments were so close to the surface and needed to further reduce their background, and because they were interested in detecting relatively light ($m \lesssim 5$ TeV) particles, they used a 15-cm-thick shield of NaI scintillator to veto all events which had an equivalent electron energy (the energy of an electron which when incident on the NaI produces the same number of electron-hole pairs as the incident h) of $E_e > 30$ keV, corresponding to $E_{\rm recoil} > 120$ keV. Thus if the cross section on I, $\sigma_{\rm I} > 120$ keV$(4)/(n_{\rm NaI} m_{\rm I} v_{\min}^2 d)$ $\sim 1.5\times10^{-22}$ cm^2 then the particles are not detected. (Here $v_{\min}^2 = 10^{-7}$ is the minimum velocity for having a significant electron-hole pair conversion efficiency, d is the thickness of the detector and the factor of 4 reflects the loss of coherence due to the large size of the I nucleus.[7] It is actually possible that h's in the high-velocity tail of the distribution would have an even larger correction for the loss of coherence, and would therefore more easily evade the veto. However, to determine this would require a detailed analysis of the iodine nucleus, one which we do not undertake.) Since 600 m.w.e. corresponds to a column density of 3.5×10^{28} GeV/cm^2, and allowing for the usual factor of 10 degradation in energy, Caldwell *et al.* set a limit of $\sigma_{\rm Earth}/m > 3\times10^{-29}$ cm^2/GeV. This bound, however, is cut off by the veto at approximately 10^6 GeV for the coherent cross sections and 10^5 GeV for the spin-dependent cross sections. This reduces the utility of their data at high masses.

The detector of Avignone *et al.* is located in the Homestake gold mine at a depth of 4000 m.w.e. Again, allowing for an energy degradation of 90% gives us a limit of $\sigma_H/m_h \gtrsim (1.6\times10^{-30}/\mathcal{F})$ cm^2/GeV. Calculation of $\mathcal{F}$ must appropriately take into account the relative abundances of the various elements in Earth's crust. Using the average composition of the crust gives $\mathcal{F}_{\rm coh} = 9\times10^5$ and $\mathcal{F}_{\rm spin} = 10^2$, although these could be somewhat altered if the composition in the vicinity of the mine differs radically from average.

We wish to determine as a function of cross section the maximum mass which each detector excludes. To do this we equate the rate of detectable events to the background over the range of allowed nuclear recoil energies (or equivalent electron or photon energies):

$$\phi\sigma N_A \frac{f}{g} = \frac{dB}{dE}\frac{2m_A v^2}{4} \ . \tag{2.4}$$

Here N_A is the number of nuclei in the detector, f is the fraction of nuclei that are "active" (unity for the coherent case, and the fraction of nuclei with spin, for the spin-dependent case), and $g=1$ for Rich, Rocchia, and Spiro and $g=2$ for Avignone *et al.* The factor of 4 on the right-hand side is the conversion factor from recoil energy to equivalent electron energy.

In the detector of Rich, Rocchia, and Spiro the background rate at equivalent-photon energies below 7 keV (Fig. 1 of Ref. 6) was approximately $dB/M\,dE = 0.1$ keV^{-1}s^{-1}g^{-1} at high altitude, and a factor of 5 below this at 100 m. The mass of the detector was 0.5 g. In this regime, so long as $E_{\rm recoil} \gtrsim 1$ keV, $E_\gamma \approx 0.25 E_{\rm recoil}$. (Below this the conversion efficiency for nuclear recoil energy into electron-hole pairs drops dramatically.) [In the overflow bin ($E_\gamma > 7$ keV), the background was $B_{\rm over} = 1$ s^{-1} at both altitudes.] Thus, the detector is sensitive to $\sigma/m_h \geq (2\times10^{-30} f)$ cm^2/GeV. (Rich, Rocchia, and Spiro claim a factor of 2 improvement on this result.) The above calculation is valid so long as one is in the long mean-free-path limit, $\sigma < \sigma_{\rm crit} = (fnd)^{-1}$, where d is the size of the detector and n is the number density of nuclei; however, once the mean free path is the length of the detector, the probability increases that the SIMP will scatter several times in the detector and therefore register in the (much higher background) overflow bin. Using Poisson statistics we can find the probability that the SIMP will scatter only once in the detector: $P(1) = ze^{-z}$, where $z = \sigma/\sigma_{\rm crit}$ is the optical depth of the detector. Condition (2.4) then applies with ϕ reduced by the factor $P(1)$. Eventually, $P(1)$ becomes small enough that the background in the overflow can compete. Then (2.4) is replaced by

$$\phi \mathcal{G}\mathcal{A}\frac{f}{g} = B_{\rm over} \ ,$$

where $\mathcal{A}$ is the detector area and $\mathcal{G}$ is a geometric factor representing the fraction of the flux which comes from a solid angle of "sky" which has a low enough column density not to degrade the h below threshold:

$$\mathcal{G} = \tfrac{1}{2}\int_y^1 \cos\theta \, d\cos\theta = \tfrac{1}{4}(1-y^2) \ ,$$

where $y = \sigma/\sigma_{\max}(m)$. Here $\sigma_{\max}(m)$ is the cross section which, for h's coming from directly overhead, causes a factor of 10 degradiation in the energy. We are thus led to a limit of $m_h < 3\times10^6 f(1-y^2)$ GeV.

For the detector of Avignone *et al.* the maximum accessible cross section as a function of mass was calculated above to be $\sigma_H/m_h \gtrsim (1.6\times10^{-30}/\mathcal{F})$ cm^2/GeV. To compute the maximum mass which can be probed as a function of cross section one must once again know the backgrounds in the detector. Here the situation is somewhat more complex than for the detector of Rich, Rocchia, and Spiro because the background is a strong function of energy; moreover, as the focus of the experiment has evolved, backgrounds have evolved differently in different regimes. The background in the detector at low E_γ ($\lesssim 36$ keV, which is the relevant range in the long mean-free-path limit) is[8] $dB/dE = 1.4\times10^{-5}$ keV^{-1} s^{-1}; hence, the detector is sensitive to $\sigma/m_h \gtrsim 2\times10^{-36}$ cm^2/GeV. Above the critical cross section $\sigma_{\rm crit}=2\times10^{-24}$ one expects $z=\sigma/\sigma_{\rm crit}$ scatterings in the detector, giving an equivalent-photon energy of $z\frac{1}{4}m_A v_0^2 = 18z$ keV. The background in the detector has been reported[9] out to 4 MeV, allowing one to probe to a z of 220. Since the detector has no overflow bin, for higher z's one must rely on the low-velocity tail of the Boltzmann distribution (this greatly dominates the spread in the energy deposition due to Poisson statistics of the number or average energy transfer of the collisions). These considerations, as well as an allowance for the geometric suppression of the flux ($\mathcal{G}$) permit us to compute the high-mass cutoff of the Avignone *et al.* results, as plotted in Figs. 1–3.

These limits from the detectors of Avignone *et al.*, Caldwell *et al.*, and Rich, Rocchia, and Spiro all appear as the bold solid lines in Figs. 1–3. In Fig. 1, of course, different portions of the curve refer to cross sections on the various substances, as described above.

B. Solid-state cosmic-ray detectors

Solid-state cosmic-ray detectors work on the same principle as the semiconductor detectors described above. One detector of interest is the main telescope of the University of Chicago instrument aboard the Pioneer 11 spacecraft. This contains seven elements D1–D7: D1–D4 and D6 are Li-drifted silicon detectors (column density of ≈ 3 mg/cm^2 each), D5 is a CsI scintillator viewed through a Li-drifted silicon detector (5.4 g/cm^2 of CsI, 1.5 mg/cm^2 of Si wafer), and D7 is a plastic scintillator crystal which fits as a sleeve around the other elements.[10] From this detector one can obtain the following limits:[11] (a) for particles which stop within the first two elements (a column density of ≈ 6 mg/cm^2) and deposit at least 3 MeV, the flux is less than 10^{-2}/cm^{2}s sr; (b) for particles which trigger (deposit at least 250 keV in each element) at least 3 elements ($\gtrsim 6$ mg/cm^2) but stop within the telescope (6 g/cm^2), the flux is less than 10^{-2}–10^{-4}/cm^{2}s sr (note that the fifth element, the CsI scintillator, has an efficiency of only 4%, so that 250 keV of detectable energy is actually 6.25 MeV of energy loss); (c) for throughgoing particles (depositing at least 250 keV in each element), the flux is less than 0.3/cm^{2}s sr.

If $m_h > 6$ MeV$/\beta^2 \approx 6$ TeV, then h has 3 MeV of kinetic energy to deposit, and can be detected in the first two elements. If the energy loss rate of h, $dE/\rho\,dx = \sigma_{\rm Si}\beta^2 > 5\times10^2$ MeV cm^2/g, then the particle will deposit 3 MeV in the first two elements; this corresponds to $\sigma_{\rm Si} > 9\times10^{-19}$ cm^2. The h will stop in the first two elements if[12]

$$\sigma_{\rm Si} \gtrsim 1.6\times10^{-19}\,{\rm cm}^2(m_h/{\rm TeV}) \times \ln\left[\frac{\sigma_{\rm Si}}{1.4\times10^{-19}\ {\rm cm}^2}\frac{\rho\Delta x_1}{{\rm mg/cm}^2}\right] .$$

If $\sigma_{\rm Si}$ is less than this, then the h will penetrate to the third element with enough kinetic energy left to produce a signal in that element. Since the minimum total energy deposition is now only 750 keV, this applies to all h's with $m_h > 750$ GeV. The third element can be triggered so long as $dE/\rho\,dx = \sigma_{\rm Si}\beta^2 > 80$ MeV cm^2/g, corresponding to $\sigma_{\rm Si} \gtrsim 1.4\times10^{-19}$ cm^2.

The h's will stop in the detector if

$$\sigma_{\rm Si} \gtrsim 1.5\times10^{-22}{\rm cm}^2\frac{m_h}{\rm TeV}\ln(10^{-6}\sigma_{\rm Si}\rho\Delta x_1/250\ {\rm keV}) .$$

Above this cross section, the flux of h's is at most 10^{-2}/cm^{2}s sr, while below it can be as large as 0.3/cm^{2}s sr. The flux of h's is approximately $10^6\beta({\rm TeV}/m_L)$cm^{2}s sr. Demanding that this be less than 10^{-2}/cm^{2}s sr (nonthroughgoing), and 0.3/cm^{2}s sr (throughgoing) gives $m_h > 10^8$ GeV and $m_h > 3\times10^6$ GeV, respectively. The region of cross section versus mass thus excluded is bordered by the light dotted line in Figs. 1–3.

C. Plastic track cosmic-ray detectors

In plastic track detectors,[13] the energy deposition from scattering of the SIMP results in the breaking of molecular bonds in the plastic; the resulting lighter molecules are etched away once the exposure is complete. If $dE/\rho\,dx \gtrsim 400$ MeV cm^2/g, then sufficient damage is caused to etch. This corresponds to $\sigma_{\rm pl} \gtrsim 7\times10^{-19}$ cm^2. Tracks must also have some minimum length, $x_{\rm min} \approx 0.25$ cm, to be seen. Following the analysis above, this implies

$$m_h \gtrsim 110\ {\rm TeV}(\rho_{\rm pl}/{\rm g\ cm}^{-3})(\sigma_{\rm pl}/10^{-19}{\rm cm}^2)(x_{\rm min}/{\rm cm}) .$$

The overall flux limits from plastic track detectors are quite severe, with the exposure being measured in m^{2}sr yr$=3\times10^{11}$ cm^2 s sr. This would be sufficient to place a limit of $m_h > 2\times10^{17}$ GeV. Even if one is skeptical of using the entire exposure to limit the flux, one can still require that it be smaller than that of the rarest cosmic rays (which are resolved into charge peaks with $\delta Z/Z \ll 1$). Taking the minimum flux to be that of fluorine[14] (3×10^{-4} cm^2 s), one finds that

$$m_h \gg 3\times10^7\ {\rm TeV} . \tag{2.5}$$

The excluded region is bordered by the light long-dashed line labeled "plastic" in Figs. 1–3.

One must be aware that plastic track detectors are only now being calibrated at $\beta \approx 10^{-3}$. However, there is no reason to expect that the energy deposition threshold at $\beta \approx 10^{-3}$ will be significantly different than at the higher values of β usually encountered.

D. Galactic-halo infall and molecular-cloud heating

If the galactic halo is to be stable, the time scale on which SIMP's lose their kinetic energy and fall into the disk must be long compared to the Hubble time. The average energy loss in the disk interstellar medium is $dE/dx = n\sigma_{tr}\Delta_E(m)$, where n is the average number density of protons (or hydrogen atoms), $n > 0.5$ cm^{-3}, and $\Delta E(m) \approx m_p v^2$ is the average energy transfer per collision (we assume $m_h > m_p$). (We are justified in using average quantities since the dynamical time of the Galaxy is much less than a Hubble time.) Since the "average" halo particle (20-kpc orbit) spends about 2% of its time in the disk, $d\ln E/dt \gtrsim 1.4\times 10^7(\sigma_{tr}/m_h)$ GeV cm^{-2}s^{-1}, where we have taken $v/c \approx 10^{-3}$. Requiring that the infall time be greater than 5×10^9 yr, implies that

$$\sigma_{tr} \le 5\times 10^{-24}\frac{m_h}{\text{GeV}}\text{cm}^2 .$$

We have neglected so far to consider the effect of shock acceleration by supernovae. It is possible that as often as once every 10^7 yr, a SIMP is shock accelerated[15] and blown either back into the halo, or out of the Galaxy. Since the binding energy of the Galaxy is $\approx 10^{60}$ ergs, while the total energy output of galactic supernovae (assuming liberally 3×10^{51} ergs per event, and a supernovae every 10 yr) is $< 3\times 10^{60}$ ergs, this possibility is marginal but not excluded.

If reinjection occurs, the infall time need only be $> 10^7$ yr, weakening the above limit by a factor of 500. Nevertheless we must take account of the heating of the disk by the dark matter. Chivukulah *et al.* considered[16] the heating of nonionized interstellar clouds by elastic collisions of the SIMP's with hydrogen, neutral or ionized. Integrating over a Maxwellian velocity distribution and assuming σ_{tr} varies slowly with velocity, they found the heating rate per nucleon to be $\gamma_h \approx 1.2 n_h \sigma_{tr} m_p v_h^3$, where $m_h \ge m_p$. Using the observed cooling rates of Pottasch, Wesselius, and van Duinen,[12] $\lambda = (8.1\pm 4.8)\ 10^{-14}$ eV/s, and requiring $\gamma_h \lesssim \lambda$ they find

$$\sigma_{tr} \lesssim 8\times 10^{-24}\ \text{cm}^2\frac{m_h}{\text{GeV}} .$$

These bounds (we use the slightly more conservative cloud cooling bound) appear as the dotted-dashed lines in Figs. 1–3.

III. MODEL-DEPENDENT LIMITS

The bounds derived above, as far as possible, depended only on the existence of a transfer cross section σ_{tr} for the SIMP on ordinary matter, and, in some cases, on the assumption that there is no significant depletion in the local halo-particle density. We can also obtain some very stringent limits on interactive dark matter given certain reasonable assumptions regarding its properties. We first discuss a bound on σ_{tr} by arguing that if there is an asymmetry either in the number density or in σ_{tr} of h's versus $\bar{h}$'s, then neutron stars would have captured enough SIMP's to form black holes in their centers and would have been absorbed by them. We next show that if SIMP's have large enough σ_{tr} to be captured by Earth (but small enough that they would have drifted to the center of Earth), then in order for neutrinos from their annihilations not to have been detected in proton-decay experiments, either the h-$\bar{h}$ asymmetry must be large or the cross section σ_{an} for annihilation into ordinary matter must be small in fact so small that, according to a standard Lee-Weinberg calculation, h's would have overclosed the Universe. In the substantial region of parameter space where both the neutron-star and Earth arguments apply, SIMP's are completely ruled out as the dark matter provided one adopts the assumptions of the Lee-Weinberg calculation. Even without the Lee-Weinberg assumptions h's with mass greater than $2\times 10^8\ m_p$ are ruled out because they would have formed a detectable black hole in the center of Earth. If σ_{tr} is large enough that h's would not yet have drifted into the center of Earth, then for σ_{an} in the cosmologically permitted range (for a symmetric abundance), more detailed analysis of the angular dependence of the signal in proton-decay experiments should either rule out these particles or detect them.

Finally, we point out that for particles which bind to ordinary nuclei, heavy-isotope searches indicate that the mass be greater than 10 TeV.

A. Longevity of neutron stars

In a previous paper[18] Gould *et al.* showed that if the dark halo consisted of charged massive particles, they would collect in the center of neutron stars, form black holes, and destroy their hosts. The existence of ~Gyr-old neutron stars placed limits on the halo density of such particles. Here we use a virtually identical argument to place upper limits on σ_{tr} for SIMP's, h, with $m_h \gg m_p$. In contrast with the case of charged dark matter, however, we assume that σ_{tr} is independent of velocity provided the latter is nonrelativistic. Actually what we obtain is a limit on $\sigma_{as} \equiv \sigma_{tr}/(p\delta_\sigma + \delta_n)^{1/p}$, where δ_σ and δ_n are the fractional particle-antiparticle asymmetries in σ_{tr} and the halo density, respectively, and $p \sim 1$ will be defined below. The limit does not apply in the case where both the number density and σ_{tr} are symmetric.

The fractional pollution of a collapsing cloud by SIMP's, h, is

$$\eta \equiv \frac{M_h}{M} = \int_{t_{\min}}^{t_{\max}} dt\ \pi[R(t)]^2\frac{\rho_h}{M} \times \int_0^{v_s(t)} d^3v\ v\sigma_{tr}F(\mathbf{v})\left[1+\frac{v_e^2}{v^2}\right] , \tag{3.1}$$

where M is the mass of the protostellar cloud, M_h is the mass of captured h, $R(t)$ is the radius of the cloud as a function of time, ρ_h is the halo density, $v_s(t)$ is the maximum velocity at infinity which can be stopped by the cloud column at time t, and $F(\mathbf{v})$ is the velocity-distribution function. We consider only those h accreted between the time when the cloud is molecular and self-gravitating ($t_{\min}$) and the time when the star is essentially formed ($t_{\max}$). We take F to be an isothermal-sphere distribution as seen by an observer moving with the disk,

$$F(\mathbf{v}) = \pi^{-3/2}v_{rot}^{-3}\exp[-(\mathbf{v}-v_{rot}\hat{\phi})^2] ,$$

where $v_{\rm rot}=220$ km s^{-1}. Proceeding as in Ref 18 we find

$$\eta(\beta)=24\sqrt{2}\pi^2\left[\frac{GM}{v_{\rm rot}^2}\right]^3\frac{\rho_h}{M}\int_{y_{\rm min}}^{y_{\rm max}}dy\, y^{-8}g(y;\beta)\ , \qquad (3.2)$$

where $y=v_e/v_{\rm rot}$, $v_e(t)$ is the escape velocity of the cloud,

$$g(y;\beta)\equiv(y^2+1.5)[2{\rm erf}(1)+{\rm erf}(z-1)-{\rm erf}(z+1)] + \pi^{-1/2}[2e^{-1}+(z-1)e^{-(z+1)^2} - (z+1)e^{-(z-1)^2}]\ ,$$

$$z(y,\beta)\equiv y(e^{2\beta y^4}-1)^{1/2}\ ,$$

and

$$\beta\equiv(v_{\rm rot}^4/4\pi G^2M)(\sigma_{\rm tr}/m_h)\ .$$

One finds[18] that $y_{\rm min}\sim0.0013$ and $y_{\rm max}\sim1.5$.

A neutron star will be destroyed[18] if tan $\gtrsim\max\{\eta_{\rm Chand},\eta_{\rm Hawk}\}$ where $\eta_{\rm Chand}\sim3.9(m_p/m_h)^2$ comes from the requirement that M_h be large enough to collapse into a black hole and $\eta_{\rm Hawk}\sim2.8\times10^{-20}$ from the requirement that the black hole be large enough to accrete neutrons. One can show that $\eta(\beta)\propto\beta^p$ where $p=1$ for ($\beta\lesssim60$), $p=0$ for $[y_{\rm min}\exp(\beta y_{\rm min}^4)\gtrsim1]$, and $p=\frac{7}{4}$ in the intervening regime. Thus $\sigma\propto m^{1-q/p}$, where $q=2$ in the Chandrasekhar regime and $q=0$ in the Hawking regime. Results of the numerical evaluation of Eq. (3.2) for $m_h<7.7\times10^7$ GeV are displayed as the short-long-dashed line in Figs. 1–3. In this region, $p=\frac{7}{4}$ and $q=2$, so that $\sigma\propto m^{-1/7}$. For $m_h>7.7\times10^7$ GeV, *already formed* neutron stars could collect enough mass in a 1×10^9 yr to form a black hole,[19] provided only that the h interact strongly enough to be stopped by a neutron star, $\sigma/m_h\gtrsim10^{-46}$ cm^2/GeV. Thus, the "neutron star" curve drops precipitously off the page at 7.7×10^7 GeV and reappears only in the lower right corner of graph.

The argument just given holds only if the h do not interact so strongly as to prevent them from drifting to the center in a reasonable amount of time. The characteristic drift time through the degenerate neutron gas is related to the cross section by[18] $\sigma/m_h\sim G\tau_{\rm drift}(m_p/2E_f)^{1/2}(E_f/kT)$, where $E_f\sim300$ MeV is the Fermi energy of the neutrons and $T\sim10^7$ K is the temperature of a neutron star after several Myr. At the high cross sections relevant to this discussion, the neutron star will be polluted with H at a fractional rate $\eta\sim0.1$, so that only $\sim\ln(5/\eta)\sim4$ drift times will be required to form a black hole. Thus, a drift time of $\lesssim100$ Myr will be short enough to destroy the neutron star. This implies that the neutron-star argument will hold provided that $\sigma/m_h\lesssim4\times10^{-21}$ cm^2/GeV. Since these cross sections lie well above the "cooling-infall" line, they are already strongly ruled out. Because of this, and because we do not want to further clutter our graphs, we do not display the neutron-star drift limit.

B. Capture and annihilation in Earth

We can use the fact that the SIMP's are efficiently captured by Earth to place significant restrictions on their masses and annihilation cross sections.[20,21,7] To do so, we restrict our attention to those particles with cross sections (with iron) which are small enough so that they will drift to the center of Earth in much less than an Earth lifetime. The drift time $\tau_{\rm drift}$ is found by balancing the gravitational force against the viscous drag, and is, in the high-mass limit,[18]

$$\tau_{\rm drift}(r)\sim\frac{12}{5}\pi^{-3/2}\frac{\sigma}{m_h}\frac{v(r)}{G}\frac{\rho(r)}{\bar\rho(r)}\ , \qquad (3.3)$$

where $\rho(r)$ and $v(r)$ are the local density and thermal speed of the iron and $\bar\rho(r)$ is the mean density interior to r. From the form of Eq. (3.3) it is clear that $\tau_{\rm drift}$ is largest near the center of Earth. To allow several e-foldings, we demand $\tau_{\rm drift}\lesssim100$ Myr, so that

$$\frac{\sigma}{m_h}\lesssim5\times10^{-21}\ {\rm cm^2\,GeV^{-1}}\ , \qquad (3.4)$$

where we have made a conservative estimate for the central temperature of Earth, ~7000 K. For masses below 100 GeV, the dependence of σ on m_h is substantially more complicated than Eq. (3.4) because the relative cross sections of various Earth materials (in both the mantle and core) are complicated functions of mass in this regime. We have taken this fact (as well as the 3-GeV evaporation mass[22] which provides the lower cutoff) into account when plotting the 100-Myr Earth-drift line (bold short dashes) in Figs. 1–3. The arguments given below apply to h's to the lower right of this line.

Provided that the cross section with iron (short dashes) is at least

$$\frac{\sigma}{m_h}\gtrsim\ln(v_h/v_{\rm esc})(f_{\rm Fe}\rho_\oplus R_\oplus)^{-1} \sim6\times10^{-33}\ {\rm cm^2\,GeV^{-1}}\ , \qquad (3.5)$$

Earth will capture virtually all particles incident upon it. Here $f_{\rm Fe}\sim\frac{1}{3}$ is the fraction of Earth which is iron and $v_h/v_{\rm esc}\sim30$ is the ratio of the SIMP speed to the escape speed from Earth. Under this condition, the captured mass of SIMP's will be

$$M_h=\pi R_\oplus^2\rho_h v_h\tau_\oplus\sim2\times10^{-15}\ M_\odot\ . \qquad (3.6)$$

If Eq. (3.4) is satisfied, these particles will drift to the center of Earth and form an isothermal distribution with effective volume[23] $V_{\rm eff}=(3T_c/2Gm_h\rho_c)^{3/2}$.

We now determine the critical-annihilation cross section $\langle\sigma_{\rm an}v\rangle_{\rm crit}$. Above this cross section, the SIMP's will reach a steady state between annihilation and capture during the lifetime of Earth. Below it, the annihilation rate will still be below the capture rate at the present time. One finds[23]

$$\langle\sigma_{\rm an}v\rangle_{\rm crit}=\sqrt{32}\frac{V_{\rm eff}}{\tau_\oplus}\frac{m_h}{M_h} \sim2\times10^{-32}(m_h/m_p)^{-1/2}\ {\rm cm^3\,s^{-1}}\ . \qquad (3.7)$$

This is an unreasonably low cross section for a particle with such high interaction cross sections. However, we will now present strong evidence that the actual annihilation cross section should be several orders of magnitude below (3.7).

In almost any particle-physics model for h, with m_h greater than a few GeV, at least some of the annihilation decay products will be ordinary Dirac neutrinos. Even if there is no branching ratio for annihilation directly into neutrinos, given the size of the interaction cross section with hadrons, it is unnatural to have no significant branching ratio into hadrons. The hadronic decay products will be highly relativistic and will interact strongly with the ambient medium, thereby generating neutrinos. We designate by ξ the fraction of decay energy in the form of electron neutrinos and antineutrinos with energies $\gtrsim$ GeV (but for simplicity, $\lesssim 10$ TeV). The neutrino energy production will then be just ξ times the mass-capture rate if the annihilation cross section is above critical, and will be scaled down by $\langle \sigma_{an} v \rangle / \langle \sigma_{an} v \rangle_{crit}$ if this cross section is below critical. We note that the cross section for neutrinos of the relevant energy range in proton-decay detectors is $\sigma_{\nu,\mathrm{nuc}} \sim \sigma_0 (E/m_p c^2)$, where $\sigma_0 = 3\times 10^{-39}$ cm^2. Taking the working life of the world's proton-decay detectors to be $m_p N_{det} \tau_{det} \sim 4$ kton yr, we predict a total neutrino detection of

$$N_\nu = \frac{N_{det}\tau_{det}\sigma_0 v_h \rho_h}{4m_p}\xi \min\left\{1, \frac{\langle \sigma_{an} v\rangle}{\langle \sigma_{an} v\rangle_{crit}}\right\}$$

$$\sim 10^9 \xi \min\left\{1, \frac{\langle \sigma_{an} v\rangle}{\langle \sigma_{an} v\rangle_{crit}}\right\} . \tag{3.8}$$

Under the assumption that $\xi \gtrsim 10^{-6}$, the failure to observe this predicted signal proves that the annihilation cross section is at least 3 orders of magnitude below critical. Alternatively, one might assume a cosmic asymmetry of $\gtrsim 99.9\%$. However, as discussed above, even a very small asymmetry would lead to the destruction of neutron stars over much of the parameter space we are considering.[24]

The annihilation cross section ($\lesssim 10^{-35}$ cm^3 s^{-1}) given by Eqs. (3.7) and (3.8) is not only unnaturally low, it is inconsistent with the results of a standard Lee-Weinberg[25,26] calculation which predicts

$$\langle \sigma_{an} v\rangle_{LW} = \frac{3\times 10^{-27}\ \mathrm{cm^3 s^{-1}}}{\Omega_h h_0^2}\left[\frac{g_f}{107}\right]^{-1/2}\frac{1}{25 x_f} , \tag{3.9}$$

where Ω_h is the fraction of critical density in h's, h_0 is the reduced Hubble constant, g_f is the number of effective relativistic degrees of freedom at freeze-out [normalized to g(photons)=2], and $x_f(\sim \frac{1}{20})$ is the ratio of the freeze-out temperature to the mass. If h's make up the dark matter and do not overclose the Universe, then $0.1 \lesssim \Omega \lesssim 2$ so that $10^{-27} \lesssim \langle \sigma_{an} v\rangle_{LW} \lesssim 10^{-25}$. It is possible to circumvent Eq. (3.9) by invoking late entropy dumping or other nonstandard mechanisms. However, if we assume Eq. (3.9) is valid, then the entire regions of parameter space above the neutron-star line and between the Earth-drift and Earth-capture lines in Figs. 2 and 3 are ruled out.

If the annihilation cross section was actually as high as (3.9), and the interaction cross sections were as high as the "Earth-drift" line, then for $m_h \lesssim 10^5$ GeV, the h's would not actually have time enough to drift to the center of Earth before they annihilated with one another. One may show, however, that they would have enough time to drift past the proton-decay detectors (at a depth ~ 1 km). In this case, the expected neutrino signal would come not from a point source at Earth's center, but from half of the "sky." That is, one would expect a cumulative excess of $10^9\xi$ events of energy $\gtrsim$ GeV coming from the "down" side of the detector. We are unaware of any studies which place limits on such an asymmetry, but we feel certain that any significant asymmetry would have been noticed.

If the interaction cross section were high enough so that the SIMP's did not drift past the proton-decay detectors before annihilating, there would still be a recognizable annihilation signal. As a simple example, we consider annihilations taking place in a spherical shell of radius r, and observed in a detector at radius $a = (1-\varepsilon)r$. By a straightforward geometrical argument one finds that the distribution of neutrino signal $N_\nu(\theta)$ would scale as

$$N_\nu(\theta) \propto [1-(1-\varepsilon)^2 \sin^2\theta]^{-1/2} . \tag{3.10}$$

The symmetry of this flux about the plane of the horizon has been noted previously in a related context.[27] The $\sec\theta$ behavior of the flux [saturating at $|\pi/2-\theta| \sim (2\varepsilon)^{1/2}$] should be easily discernible from background if it were looked for in existing data. The specific form of Eq. (3.10) is appropriate only to an underground detector in a relatively flat region. If the detector were inside a mountain, for example, the expected distribution would be quite different. However, it would still be easy to predict this distribution on the basis of local geography. We therefore believe that existing data could be used to observe or rule out symmetric, very high cross-section dark matter. In the region not excluded by the neutron-star argument, the possibility would remain that the dark matter has a cosmic asymmetry.

The Earth argument can be used to rule out a large region of parameter space even if one ignores the Lee-Weinberg constraints. For $m_h \gtrsim 2\times 10^8\ m_p$, the particles collecting at the center of Earth will attain critical mass and initiate collapse during an Earth lifetime.[18] Their collective mass will exceed the Chandraskehar mass and they will collapse to a black hole in a time $\sim \tau_{drift}/3$. During the collapse process, annihilation will accelerate, but the fraction of particles annihilated,

$$\frac{\Delta M_h}{M_h} \sim \frac{M_h}{m_h}\int dt \frac{\langle \sigma_{an} v\rangle}{\sqrt{32}V_{eff}}\frac{r_0^3}{r(t)^3}$$

$$= \frac{\tau_{drift}}{\tau_\oplus}\frac{\langle \sigma_{an} v\rangle}{\langle \sigma_{an} v\rangle_{crit}} \ln\Lambda , \tag{3.11}$$

is small. Here, $r_0 \sim V_{\rm eff}^{1/3}$ is the radius of the particle cloud when it goes critical, $\ln\Lambda = \ln(r_0 c^2/GM_h) \lesssim 50$, and we have used the fact[18] that $dr/dt \sim r^{-2}$. Once the black hole formed it would grow at the Eddington limited rate, with luminosity $L/L_\odot \sim 6.5\times 10^3\, M_{\rm bh}/M_\odot$. Assuming a conversion efficiency $\varepsilon \lesssim 10\%$, this implies an e-folding time $\lesssim 200$ Myr. At the Eddington luminosity, the flux reaching Earth's surface from the black hole would equal that from the Sun when $M_{\rm bh} \sim 7\times 10^{-14}\, M_\odot$. Since this limit would be exceeded in $\lesssim 1$ Gyr, and since the observed energy generated by Earth is of order 10^{-4} that received by the Sun,[28] all particles satisfying Eqs. (3.4) and (3.5) with masses greater than $2\times 10^8\, m_p$ are ruled out. This limit is indicated by the bold vertical dashed line in Figs. 1–3.

C. Heavy isotopes

If h's bind to nuclei, as they would, for example, if they were charged, or possibly if they carried ordinary color[29] (of course, they would be color singlets, since color is confined, but they would still feel color van der Waals forces), then they would show up in heavy isotope searches. Smith *et al.*[30] and a group at the University of Rochester have shown that the abundance of heavy hydrogen with mass less than 10 TeV is less than one part in 10^{24} by number. Nitz *et al.*[31] have searched for such objects in an electronic beam line. For $Z=3-9$ they set abundance limits of $\lesssim 10^{-20}$.

The flux at Earth of h's from the galactic halo is naively

$$\phi_h \approx 1.2\times 10^4 \frac{\rm TeV}{m_h} {\rm cm}^{-2}\,{\rm s}^{-1} . \tag{3.12}$$

If $\sigma_{\rm tr} \gtrsim 10^{-28}$ cm^2/GeV, then h's stop in the ocean. If they can bind to oxygen nuclei, they will do so at least 33% of the time, being incorporated into heavy water. (If $\sigma_{\rm tr} \gtrsim 10^{-26}$ cm^2/GeV, the h stop in the atomosphere and bind to oxygen about 25% of the time, then drift down onto the surface of Earth.) Once in water, the interaction cross section with matter is enhanced, and the drift velocity of the heavy-water molecule is $v_{\rm drift} \lesssim 2\times 10^{-4}(m_h/{\rm TeV})$km/yr. Given that the characteristic mixing time for the ocean is about 10 yr, for $m_h \lesssim 10^3$ TeV, currents will keep the ocean well mixed. Assuming a perfectly mixed ocean of 10 km uniform depth, the expected number density of hHO compared to that of H_2O is

$$\frac{n_h}{n_{\rm H_2O}} \approx 8\times 10^{-18} \frac{\rm TeV}{m_h} \frac{t_{\rm acc}}{\rm yr} . \tag{3.13}$$

Here $t_{\rm acc}$ is the time period over which hHO accumulates in the ocean and is not removed (by chemical or other processes).

For the hHO not to have been detected, it would have had to have been removed from the ocean with $t_{\rm acc} \lesssim 40$ s. Thus, if h binds to nuclei, $m_h > 10$ TeV. This limit is indicated by the vertical dotted-dashed line in Figs. 1–3.

IV. CONCLUSION

A. The allowed regions

We have seen that limits can be imposed on the cross section of the dark matter from many sources. As stated above, in Fig. 1 we present the raw cross-section limits, making no attempt to relate cross sections off nuclei of different A and Z. This we do for the use of those who may invent models in which this relationship is nonstandard.

In Figs. 2 and 3 we plot the limits on the dark-matter cross section on protons as a function of mass for coherent and spin-dependent cross sections, respectively. We have taken $\sigma(A,Z) = A^2[m_{\rm red}(A)/m_{\rm red}(1)]^2\sigma_p$ for the coherent case and $\sigma(A,Z) = ({\rm spin})[m_{\rm red}(A)/m_{\rm red}(1)]^2\sigma_p$ for the case of spin-dependent cross sections. That is, A^2 from coherence and $[m_{\rm red}(A)/m_{\rm red}(1)]^2$ from phase space. All SIMP's lying within the region enclosed by the bold solid line labeled "balloon/mine" are excluded, as are SIMP's which lie above the light dash-dotted line labeled "cooling/infall," or inside the regions enclosed by the dotted line labeled "Pioneer 11" or the dashed line labeled "plastic." The region between the heavy dashed line marked "Earth drift" and the light dashed line marked "Earth capture," and to the right of the heavy dashed line marked "Earth black hole" is also excluded. The region between the "Earth drift" line and the "Earth capture" line, and to the left of the "Earth black hole" line is also excluded unless there is a large enough asymmetry in the terrestial abundance of SIMP's and anti-SIMP's to eliminate any observable neutrino signal arising from SIMP–anti-SIMP annihilations (as described in Sec. III B). (The terrestial asymmetry can be due to an asymmetry in the halo abundance of SIMP's and anti-SIMP's, or to an asymmetry in their cross sections on matter, hence the probability of their capture by Earth.) (For this last region, it is also possible that the SIMP's have unnaturally small annihilation cross sections, and that their abundance has been diluted after freeze-out by a late-time entropy dumping or inflation.) The region lying above the short-long-dashed line marked "neutron star" (note the unmarked continuation of this line in the lower right-hand corner of each graph) is excluded so long as there is *neither* a significant asymmetry in the SIMP versus anti-SMP abundances *nor* one in their transfer cross sections on matter. The region to the left of the dotted-dashed line marked "heavy isotope" is excluded if the SIMP's bind to light nuclei (H,Li,C,N,O).

The remaining allowed regions have been shaded for identification. The regions shaded with vertical hatch are allowed only if there is a SIMP–anti-SIMP asymmetry sufficient to suppress the neutrino signal from annihilations in Earth. The regions shaded with horizontal hatch are allowed so long as there is not a significant SIMP–anti-SIMP asymmetry (so that the neutron-star limit applies). There are no restrictions on the cross-hatched regions. It is notable that there are regions of each type in which the cross sections off nuclei are less than geometric.

[In the region to the left of the "Earth black-hole" line, the shading reflects the assumption that the SIMP's have annihilation cross sections which obey the Lee-Weinberg condition (3.9). If we relax this assumption, and invoke late time entropy dumping to dilute the SIMP abundance, and if the annihilation cross section is $\gtrsim 10$ orders of magnitude below the Lee-Weinberg value, then the bounds from neutrino production in Earth would be evaded. The effect of this on the shading of the figures is to add horizontal hatching to the region between the "Earth drift," "balloon mine," and "Earth black-hole" lines, except for the regions above the "cooling/infall," "Pioneer 11," and "plastic" lines.]

B. Filling the windows

The window covering the most conventional region of parameter space is the one lying above the "balloon/mine" line and below 10^5 GeV. Within this region the cross section of SIMP off nitrogen (or oxygen or silicon) must be at least 10 barns, which is ten times bigger than the nuclear geometric cross section. Therefore the scattering must be mediated by a light particle which can give rise to a long-range force, or there must exist bound-states resonances. The former possibility has two severe difficulties. First, the contribution of this exchange force to the cross sections of protons and neutrons off each other must not overwhelm the ordinary strong-interaction cross sections. Also the exchange particle is quite light ($\approx$MeV) and couples strongly to ordinary matter, yet has not been seen in colliders. We doubt the possibility of overcoming these two objections.

The possibility of resonances is also problematic. If there is a bound state of the SIMP with light nuclei, then the mass of the SIMP must be >10 TeV. From Fig. 1 we see that for these masses, the SIMP must have a cross section on nitrogen which is more than 10^4 times the geometric cross section; moreover, the cross section on protons must be at least a factor of 10 smaller than this. One could also argue, however, that it is at least possible both to use the resonance mechanism and to evade the heavy isotope bounds. Such would be the case if the resonances were due not to bound states but to "quasibound states." (This effect explains the large cross section for low-energy neutron-neutron scattering.[32])

The allowed region between 10^5 and 10^8 GeV contains a small subregion (at the end of the finger) which for the coherent case has less than geometric cross section (and evades another oft-discussed bound that the mass of a fundamental dark matter particle must be less than 350 TeV in the absence of late-time entropy dumping). Apart from this subregion, this region suffers from the same problems as discussed above. As discussed below, this region is the most readily probed by further experiments.

Much of the allowed above 10^{10} GeV and below the "Earth capture" line contains cross sections which are less than geometric. In addition, the argument that late-time entropy dumping must be avoided becomes progressively less compelling as m_h approachs the grand-unified-theory (GUT) mass. However, individual models would need to ensure that the exchange particle mediating the SIMP-matter scattering would not have been seen in p-$\bar{p}$ colliders.

The allowed region above 10^{10} GeV, above the "Earth drift" line and below the "cooling infall" line would appear to be of little interest as the cross sections are absurdly high from a particle-physics standpoint.

C. New experiments

Several of the allowed regions could be probed without significant technological advances. By far the simplest undertaking would be a reanalysis of the proton-decay data to search for a "surface signal" as described in Sec. III. A negative result would raise the Earth drift line above the "cooling infall" line. A second experiment which is also relatively straightforward is to place one of the solid-state "mine" or "balloon" detectors at Earth's surface and surround it by iron. It is important that the shielding be Fe (although Cu would be marginally acceptable) as opposed to Pb so that the degradiation of the SIMP energy can be completely understood. Similarly, if the experiment is performed indoors, the composition and geometry of the building should be known, and there should not be a significant abundance of massive elements ($A \gtrsim 56$). The detector readout should contain an overflow (i.e., high-energy) bin. The amount of shielding should be progressively increased until the background in the overflow bin is $\lesssim 10^{-2}$ g^{-1}s^{-1}. The detector need only be run at each shielding thickness long enough to become background limited. A negative result would effectively close the window at intermediate SIMP masses, in particular eliminating the region with cross sections less than geometric. (The latter objective could also be attained by a detector located as deep as 600 m.w.e. below Earth's surface, but in the absence of any cosmic-ray veto or heavy-element shielding.)

Some advances could be made in the high-mass and low-cross-section regime simply by putting an overflow bin on existing underground detectors, or by installing a high-energy ($E \gtrsim 5$ MeV) detector alongside them.

Finally, the "balloon/mine" line could be pushed to higher cross sections at low masses by placing the balloon-type detector in a rocket, with as thin a window as possible between the detector and space.

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$$\sigma_{\mathrm{Si}} \gtrsim 10^{-21}\ \mathrm{cm}^2 \frac{\mathrm{g/cm}^2}{\rho\Delta x}\frac{m_h}{\mathrm{TeV}} \ln(10^{-6}\sigma_{\mathrm{Si}\rho}\Delta x_1/250\ keV)\ .$$

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