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EXPERIMENTAL CONSTRAINTS ON DARK MATTER

JOHN ELLIS

CERN - Geneva

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1. INTRODUCTION

In this talk I will discuss some of the particle candidates for dark matter that I personally find the most plausible. For each one, I will review the constraints on it that come from both accelerator and non-accelerator experiments. For discussions of the astrophysical evidence for dark matter see [1]. Other recent reviews on dark matter include [2]. Among the new topics reviewed here are the constraints on dark matter candidates given by LEP [3], and a recent resurgence of interest [4] in metastable particles from a hidden sector of string theory as possible candidates for dark matter.

2. MASSIVE NEUTRINOS

Neutrinos have one unique virtue among all the particles that I discuss: they are already known to exist. It is also well known that a neutrino can have the critical density ρ_c suggested by inflation if its mass falls into either of two ranges: $m_\nu \sim 10$ to 100 eV or $m_\nu \sim 1.3$ to 7 GeV [5]. The former range occurs when the relic neutrino density is almost independent of mass, and hence the mass density ρ_ν is almost linear in m_ν :

$$\frac{\rho_\nu}{\rho_c} \simeq \Sigma_\nu \left(\frac{m_\nu}{100 \text{ eV}} \right) h_0^{-2} \quad (1)$$

where h_0 is the present Hubble expansion rate in units of $100 \text{ kms}^{-1} \text{ Mpc}^{-1}$. There is then an unacceptable intermediate range of m_ν where $\rho_\nu > \rho_c$, but then ρ_ν falls again, suppressed by a Boltzmann factor, yielding the second allowed range. Less well known is a third range of m_ν where $\rho_\nu \sim \rho_c$ is possible, namely $m_\nu \sim 100 \text{ GeV}$ to 1 TeV [6]. Beyond this range, ρ_ν is again greater than ρ_c because the maximum conceivable annihilation cross-section $\sigma \sim 1/m^2$ falls, leaving behind too many relic neutrinos*.

If the neutrino mass falls into the lightest range, then neutrinos would constitute hot dark matter. This mass range is certainly allowed for the μ and τ neutrinos, for which laboratory experiments tell us [7]:

$$m_{\nu_\mu} < 250 \text{ keV}, \quad m_{\nu_\tau} < 35 \text{ MeV} \quad (2)$$

but now seems to be experimentally excluded for the electron neutrino:

$$\begin{aligned} m_{\nu_e} &< 15.4 \text{ eV (Zürich [8])} \\ &9.4 \text{ eV (Los Alamos[9])} \\ &8 \text{ eV (World data[10])} \end{aligned} \quad (3)$$

* There is a similar pattern of allowed mass ranges for other particles, as we shall see later.

whereas the SN 1987a neutrino data do not exclude $m_{\nu_e} \sim 20 \text{ eV}$ [7].

An interesting new development is the preliminary report from the SAGE solar neutrino experiment [11], which sees no neutrino signal in its first few months of data-taking, compared with an expected 14 ± 2 events. This experiment is sensitive to the low-energy pp neutrinos which are produced by the Sun's major energy-production mechanism, whose rate can be predicted reliably [12]. Therefore, this discrepancy, if confirmed by further data-taking, could only mean that some intrinsic neutrino property causes flavour mixing between the Sun and the Earth. Candidates are a neutrino magnetic moment and neutrino masses. In the former case [13], one might expect time variations in the rate of neutrino detection in terrestrial experiments correlated with the sunspot cycle and the season of the year. Such an effect has been claimed in the Homestake mine data, but is not seen in the Kamiokande solar neutrino data [14]. Perhaps neutrino masses are non-zero [15], after all ...?

Could there be more than the three known light species of neutrino? LEP tells us not, by measuring the width and height of the Z^0 peak [3]. The width of the Z^0 peak is determined, via the uncertainty principle, by its decay rate Γ_Z , which may be well approximated by

$$\Gamma_Z = \Gamma_{\text{inv}} + 3\Gamma_{\ell\ell} + \Gamma_{\text{had}} \quad (4)$$

where the e^+e^- , $\mu^+\mu^-$ and $\tau^+\tau^-$ partial widths are equal and denoted by $\Gamma_{\ell\ell}$. This and the hadronic partial width Γ_{had} can be calculated reliably in the Standard Model [16], and measurements [3] agree with these predictions. One can parametrize the invisible width Γ_{inv} as

$$\Gamma_{\text{inv}} = N_\nu \Gamma_{\nu\nu} \quad (5)$$

where in the standard model $\Gamma_{\nu\nu} \simeq 165 \text{ MeV}$ for each light neutrino species[16]. However, Γ_{inv} could also include Z^0 decays into other neutral particles, such as massive fermions, scalars, or ...? The hadronic cross-section on the Z^0 peak is given by

$$\sigma_{\text{had}}^{\text{peak}} = \frac{12\pi}{m_Z^2} \left(\frac{\Gamma_{\ell\ell}}{\Gamma_Z} \right) \left(\frac{\Gamma_{\text{had}}}{\Gamma_Z} \right) (1 - \delta_{RC}) \quad (6)$$

where δ_{RC} is a radiative correction. Because of the Γ_Z factors in the denominator of (6), and the dependence ((4), (5)) of Γ_Z on N_ν , $\sigma_{\text{had}}^{\text{peak}}$ would change by 13% for $\Delta N_\nu = 1$. By measuring this peak Z^0 cross-section and the Z^0 width, LEP experiments find [3]

$$N_\nu = 2.95 \pm 0.11 \quad (7)$$

Clearly there are only three light neutrino species, and not much scope for other invisible Z^0 decay products. These must either be massive, so that Z^0 decays into them

are suppressed by phase space [17], and/or coupled to the Z^0 more weakly than conventional neutrinos.

In particular, an extra massive Dirac neutrino ν_D could only be tolerated if

$$m_{\nu_D} > 44 \text{ GeV} \quad (8)$$

whilst a massive Majorana neutrino ν_M should have

$$m_{\nu_M} > 38 \text{ GeV} \quad (9)$$

and three degenerate sneutrinos $\tilde{\nu}$ would need to have

$$m_{\tilde{\nu}} > 42 \text{ GeV} \quad (10)$$

These limits mean that none of these particles could be the dark matter of the Universe unless they weigh considerably more than $M_{Z^0}/2$ [6]. For $m_\nu < M_{Z^0}/2$, from the bound (8) one infers

$$\Omega_{\nu_D} h_0^2 < 2 \times 10^{-5} \quad (11)$$

and from (9) one infers

$$\Omega_{\nu_M} h_0^2 < 2 \times 10^{-3} \quad (12)$$

whilst the range (10) is inconsistent with the underground ${}^{76}\text{Ge}$ experiments [18] discussed later.

However, these results only apply to massive neutral particles with the same couplings as the Z^0 as the conventional light neutrinos. Particles with suppressed couplings to the Z^0 , such as mixtures of conventional $SU(2)$ doublet neutrinos with singlets, could have larger relic densities ρ . These vary inversely with the annihilation cross-section $\rho \sim \frac{1}{\sigma_{\text{ann}}}$, and for a mixing angle $\phi_Z, \sigma_{\text{ann}} \propto \sin^2 \phi_Z$ so that ρ increases. One can consider two cases for such “partial neutrinos”: Majorana states ($\nu_M = \bar{\nu}_M$) for which LEP imposes [19]

$$\sin^2 \phi_Z < \frac{(N_\nu - 3)}{\beta(3 + \beta^2)/4} \quad (13a)$$

as seen in Fig. 1a, and Dirac states ($\nu_D \neq \bar{\nu}_D$), for which LEP imposes [19]

$$\sin^2 \phi_Z < \frac{(N_\nu - 3)}{\beta^3} \quad (13b)$$

as seen in Fig. 1b. Also shown in Figs. 1 are the values of $\sin^2 \phi_Z$ which make ρ_ν critical for different choices of the present Hubble expansion rate. Also shown in Fig.

1b is the ${}^{76}\text{Ge}$ constraint [18] on Dirac particles which constitute the galactic halo. We see that LEP imposes [19]

$$m_{\nu_\mu} \gtrsim 19 \text{ GeV (cf. 7 GeV previously)} \quad (14a)$$

for $h_0 = 1/2$, and

$$M_{\nu_D} \gtrsim 19 \text{ GeV (cf. 4 GeV previously)} \quad (14b)$$

except for a small hole vulnerable to future ${}^{76}\text{Ge}$ experiments.

These are just some examples how LEP requires dark matter candidates to be heavier than previously considered: we will see more later.

3. AXIONS

The axion is a light pseudoscalar boson introduced by the Peccei-Quinn [20] solution for conserving CP in the strong interactions. The axion mass and couplings

$$m_a \sim \frac{(100 \text{ MeV})^2}{f_a}, \quad g_{a\text{ff}} \sim \frac{m_f}{f_a} g_{a\tau\tau} \sim \frac{\alpha}{\pi f_a} \quad (15)$$

scale inversely with the axion decay constant f_a that is often related to some large Higgs v.e.v. beyond the Standard Model. It had been thought that the dominant source of axion energy density in the present-day Universe would be coherent waves that gave non-relativistic cold dark matter [21]. The density of these waves would give $\Omega_a \leq 1$ if

$$f_a \lesssim 10^{12} \text{ GeV} \quad (16)$$

in the absence of inflation, whilst larger values of f_a could be tolerated if we were at a special place in an inflationary Universe [22]. Recently, it has been suggested that axion strings might radiate sufficiently many axions to require strengthening the bound (16) to $f_a \lesssim 10^{10} \text{ GeV}$ [23]. (Although emitted relativistically, these axions would now also be non-relativistic.) However, this strengthening of the bound (16) is still open to question [24].

Many laboratory experiments combine to tell us that $f_a \gtrsim 10^3 \text{ GeV}$, so that the axion must arise from physics beyond the Standard Model. The Standard Solar Model and searches [25] for the solar axio-electric effect independently tell us that $f_a \gtrsim 10^7 \text{ GeV}$. Other astrophysical objects such as Red Giants and White Dwarfs increase this lower limit to $f_a \gtrsim 10^9$ to 10^{10} GeV [26], leaving a small window below the cosmological upper limit (16).

This window was partly closed by the observation [27] of a neutrino pulse from Supernova 1987a with a characteristic temperature of order 3 to 5 MeV and total energy output $\sim (2 \text{ to } 4) \times 10^{53} \text{ ergs}$, as expected from collapse calculations [28] in which 99% of the remnant neutron star's binding energy is emitted via neutrinos. These observations limited the energy available to be radiated via axions to $\lesssim 3 \times 10^{53} \text{ ergs}$. Conventional axions with $f_a \gtrsim 10^8 \text{ GeV}$ as indicated by the previous astrophysical limits would escape freely from the core of the neutron star where the density rises to $\sim 10^{15} \text{ gm}^{-3}$ and the temperature to as much as 70 MeV. The dominant axion emission process is expected to be the bremsstrahlung reaction $N + N \rightarrow N + N + a$, which is consistent with the neutrino data if [29]

$$f_a \gtrsim 3 \times 10^{10} \text{ GeV} \quad (17)$$

with some uncertainty depending on the model for axion-nucleon couplings as seen in Fig. 2 (using the naive quark model or the EMC data on the spin content of the proton discussed later) and on the nuclear equation of state [30].

The latter could have a big effect, particularly if the neutron star core is actually made of quark-gluon plasma [31], which is not generally expected but cannot be ruled out. The dominant axion emission processes would then be $g + q \rightarrow q + a$, $q + \bar{q} \rightarrow g + a$, and transverse $g \rightarrow \text{longitudinal } g + a$ decay. For a star with mass $\sim 1.4 M_\odot$ and core density $\sim 10^{15} \text{ gcm}^{-3}$, the quark chemical potential would be ~ 200 to 500 MeV and the resulting lower bound on the axion decay constant would be [31]

$$f_a \gtrsim 2 \times 10^8 \text{ GeV} \quad (18)$$

as seen in Fig. 3, two orders of magnitude weaker than the conventional estimate (17). Because of this, laboratory experiments should try to explore all the window between $f_a \sim 10^{10} \text{ GeV}$ and $\sim 10^{12} \text{ GeV}$.

One type of experiment has been proposed by Sikivie [32] that might be sensitive in this range. One looks for axion conversion to a microwave photon in a strong magnetic field. For $f_a \sim 10^{12} \text{ GeV}$, the axion mass corresponds to a frequency $\nu \sim 3 \text{ GHz}$ and the typical power in a well-tuned cavity would be $\sim 10^{-19} \text{ W}$ if axions constitute the galactic halo. Frequencies between ~ 1 and $\sim 4 \text{ GHz}$ have been scanned [33], but with a sensitivity several orders of magnitude weaker than expected if axions constitute the galactic halo, as seen in Fig. 4. However, there are prospects for improving the sensitivity sufficiently to exclude an axion halo with $f_a \sim 10^{12} \text{ GeV}$.

4. LIGHTTEST SUPERSYMMETRIC PARTICLE

Most supersymmetric theories contain one stable particle which should be around today as a cosmological relic. This is because of a multiplicatively conserved quantum number called R parity [34] which is related to baryon number B , lepton number L and spin S :

$$R = (-1)^{3B+L+2S} \quad (19)$$

Clearly any theory conserving B and L must also conserve R , which takes the values $+1$ for all particles in the Standard Model, and -1 for all their conjectured spartners. Even GUTs which violate B and L separately tend to conserve R , but L and hence R could in principle be violated by a v.e.v. for a sneutrino field [35] or an explicit bilinear Higgs-lepton coupling in the supersymmetric Lagrangian [36]. However, these possibilities are severely constrained by upper limits on $\Delta L \neq 0$ processes. In the presence of R conservation, (1) sparticles are always produced in pairs, e.g., $e^+e^- \rightarrow \tilde{\mu}^+\tilde{\mu}^-$ or $\tilde{p}\tilde{p} \rightarrow \tilde{q}\tilde{q} + X$, (2) heavier sparticles decay into lighter sparticles, e.g., $\tilde{e} \rightarrow e\tilde{\gamma}$ and $\tilde{q} \rightarrow q\tilde{g}$, and (3) the lightest sparticle is stable because it has no allowed decay mode.

If such a relic particle had either electromagnetic charge or strong interactions, it would have condensed on to Earth along with ordinary matter, and would be detectable as an anomalous heavy isotope with a calculable abundance [37].

$$\frac{n(\text{relic})}{n(\text{proton})} \propto \frac{1}{\sigma_{\text{ann}}} \sim \begin{cases} 10^{-10} & (\text{strong interactions}) \\ 10^{-6} & (\text{electroweak interactions}) \end{cases} \quad (20)$$

This estimate conflicts with the experimental upper limits [38]

$$\frac{n(\text{relic})}{n(\text{proton})} \lesssim 10^{-15} \text{ to } 10^{-30} \quad (21)$$

obtained for $1 \text{ GeV} \lesssim m_{\text{relic}} \lesssim 10^4 \text{ GeV}$ as shown in Fig. 5. Hence any stable relic particle must be electromagnetically neutral and possess only weak interactions. The possible supersymmetric candidates are sneutrinos $\tilde{\nu}$ of spin 0, some photino $\tilde{\gamma}/\text{higgsino } \tilde{h}/\text{zino } \tilde{Z}$ combination of spin $1/2$, known generically as a neutralino, and the gravitino $\tilde{G}$ of spin $3/2$.

Sneutrinos would have coherent interactions with heavy nuclei that are as strong as those of Dirac neutrinos, and hence the underground ${}^{76}\text{Ge}$ experiments [18] also tell us they cannot constitute the galactic halo if $10 \text{ GeV} \lesssim m_{\tilde{\nu}} \lesssim 1 \text{ TeV}$. Moreover, galactic halo electron or muon sneutrinos weighing more than about 3 GeV would have been

captured by the Sun and then annihilated to give high-energy solar neutrinos that have been excluded by proton decay experiments [39]. Sneutrinos lighter than 42 GeV are excluded by LEP (10) [3], so we conclude that the halo cannot be made of sneutrinos.

Gravitinos could close the Universe in the absence of inflation if $m_{\tilde{G}} \sim 1$ keV, but models suggest that $m_{\tilde{G}} > m_{\tilde{t}} \gtrsim 40$ GeV. Hence gravitinos seem unlikely to constitute the dark matter, but this cannot be excluded. Gravitinos only have very weak interactions with ordinary matter, and hence are likely to be undetectable, so we should hope that they are not the dark matter!

Photinos mix in general with other spin-1/2 particles such as Higgsinos, and both are plausible candidates for the dark matter in supersymmetric models [40]. Their mass mixing matrix in the minimal supersymmetric extension of the Standard Model takes the form

$$\begin{pmatrix} M_2 & 0 & \frac{-g_2 v_2}{\sqrt{2}} & \frac{g_2 v_1}{\sqrt{2}} \\ 0 & M_1 & \frac{g_1 v_2}{\sqrt{2}} & -\frac{g_1 v_1}{\sqrt{2}} \\ \frac{-g_2 v_2}{\sqrt{2}} & \frac{g_1 v_2}{\sqrt{2}} & 0 & \mu \\ \frac{g_2 v_1}{\sqrt{2}} & -\frac{g_1 v_1}{\sqrt{2}} & \mu & 0 \end{pmatrix} \quad (22)$$

in the $\tilde{W}^3/\tilde{B}/\tilde{H}_2^0/\tilde{H}_1^0$ basis, where $M_{2,1}$ are gaugino masses commonly expected to be in the ratio of the corresponding electroweak couplings (α_2, α_1) , $v_{1,2} \equiv \langle 0|H_{1,2}^0|0\rangle$, and μ is an $H_1 H_2$ mixing term. The lightest neutralino (LSP: denoted by χ) state is a photino in the limit $M_{1,2} \rightarrow 0$:

$$m_{\tilde{\gamma}} \simeq \frac{8}{3} \frac{g'^2}{g'^2 + g_2^2} M_2 \quad (23)$$

and a higgsino in the limit $\mu \rightarrow 0$. The same parameters also control the masses of the chargino partners of the $W^\pm$ and $H^\pm$. In most supersymmetric models [41] one finds

$$m_{\tilde{q}} \gtrsim m_{\tilde{t}} \simeq m_{\tilde{b}} \gg m_\chi \quad (24)$$

so that the LSPs χ are the most plausible candidates for the dark matter.

Moreover, in contrast to many other dark matter candidates, the lightest neutralino naturally has a relic density close to the critical one. As for neutrinos, the relic neutralino density $\propto 1/\sigma_{ann}$, while the relic annihilation diagrams involve sfermion $\tilde{f}$, Z^0 and Higgs exchanges. Their rates are controlled by $m_{\tilde{f}}, \tan\beta \equiv v_2/v_1$, and the Higgs masses that may be parametrized by the pseudoscalar mass m_A . Plausible values of $(m_{\tilde{f}}, \tan\beta, m_A)$ give a relic neutralino density close to the critical density: $\Omega_\chi \simeq 1$, and the accelerator constraints on sparticle masses do not allow $\Omega_\chi \ll 1$ [42].

The following are the principal experimental constraints on neutralinos and the related $\tilde{W}^\pm/\tilde{H}^\pm$ states $\chi^\pm$ known as charginos. (1) LEP tells us that $m_{\chi^\pm} > 45$ GeV [43]. (2) LEP tells us that extra supersymmetric decay modes of the Z must have $\Delta\Gamma_Z < 175$ MeV [3]. (3) Using this, UA2 and CDF measurements [44] of $\sigma_W B(W \rightarrow e\nu)/\sigma_Z B(Z \rightarrow e^+e^-)$ constrain supersymmetric decay modes of the W . (4) LEP tells us (7) that invisible decay modes of the Z^0 correspond to 2.95 ± 0.11 equivalent neutrino species [3]. (5) LEP provides coupled constraints on the cross-section at the Z^0 peak and the total Z^0 decay width [3]. (6) UA2 and CDF tell us that $m_{\tilde{g}} \gtrsim 79$ GeV [44]. (7) ALEPH [45] has searched unsuccessfully for “zen” events due to $Z^0 \rightarrow \chi\chi', \chi' \rightarrow \chi + \text{visible particles}$ where χ' is a heavier neutralino. Fig. 6 shows the impacts [46] of all these constraints on the (μ, M_2) parameter plane for the choice $\tan\beta = 2$. Figure 7 shows the domains of the (μ, M_2) parameter space where $\Omega_\chi = 1$ can be obtained for some choice of $(m_{\tilde{f}}, \tan\beta, m_A)$ [47].

It is possible to distinguish four characteristic regions where $1/4 < \Omega h_0^2 < 1/2$ is possible, and we have explored [19] the values of Ω_χ that can be obtained for different choices of these parameters when (μ, M_2) lie along each of the four rays shown in Fig. 7. (1) M_2 small, μ large and negative: $\mu/M_2 \simeq -5$. In this region the LSP is mainly a gaugino ($\tilde{\gamma}$ or $\tilde{B}$), the preferred range of m_χ increases with $m_{\tilde{q}}, Z^0$ exchange is not very important, but H exchange can be important if $m_H \sim 2m_\chi$. Values of Ωh_0^2 as functions of m_χ for different choices of $(m_{\tilde{q}}, \tan\beta, m_A)$ along the ray (a) are shown in Fig. 8 [19]. We see that $\Omega h_0^2 \lesssim 1/2$ only if $m_\chi \gtrsim 20$ GeV, except in exceptional cases where $m_\chi \sim \frac{1}{2}m_H : m_H < 40$ GeV. If $m_{\tilde{q}} = 10^3$ GeV, we see that $\frac{1}{4} \lesssim \Omega h_0^2 \lesssim \frac{1}{2}$ only for $m_\chi \sim 50$ to 60 GeV [19]. (2) M_2 large, μ small and negative: $\mu/M_2 \simeq -1/10$. In this region the LSP is mainly a Higgsino, its annihilation is dominated by Z^0 exchange and relatively insensitive to $\tilde{q}$ and H exchange, and $\frac{1}{4} \lesssim \Omega h_0^2 \lesssim \frac{1}{2}$ only for $m_\chi \gtrsim 60$ GeV [19]. (3) M_2 large, μ small and positive: $\mu/M_2 \simeq 1/10$. This case is similar to the previous one, with $\frac{1}{4} \lesssim \Omega h_0^2 \lesssim \frac{1}{2}$ only for $m_\chi \gtrsim 50$ GeV [19]. (4) M_2 small, μ large and positive: $\mu/M_2 \simeq 5$. In this case the LSP is again mainly a gaugino, the preferred range of m_χ increases with $m_{\tilde{q}}$. However, Z^0 exchange is more important, and H exchange less important, than in case (1) above. In this region we always find $m_\chi \gtrsim 30$ GeV. Values of Ω_χ^2 along the rays (b), (c) (d) for different choices of $(m_{\tilde{q}}, \tan\beta, m_A)$ are shown in [19].

Thus we see that although LEP (and other experiments) strongly constrain neutralino candidates for the dark matter, they do not exclude it (unlike the massive neutrino and sneutrino cases discussed earlier). We will now explore strategies for detecting neutralino dark matter.

5. POSSIBLE WAYS TO DETECT SUPERSYMMETRIC RELICS

5.1 ANNIHILATION IN THE GALACTIC HALO

If the galactic halo is largely composed of supersymmetric relics, pairs of them will occasionally meet and annihilate into conventional fermion-antifermion pairs ($\ell\bar{\ell}$, $q\bar{q}$). These will in turn produce [47] fluxes of stable particles such as $\bar{p}$, e^+ , γ and ν that can in principle be detected amongst the cosmic rays. A typical stable particle flux

$$\phi_{\bar{p}} \text{ or } e^+ \text{ or } \gamma \text{ or } \nu \propto n^2 \cdot \sigma_{\text{ann}} \cdot \tau \quad (25)$$

where n is the number density of relic particles in the halo, σ_{ann} is their annihilation cross-section, and τ is the length of time the stable particle spends in the galactic halo. The latter is easy to calculate for photons or neutrinos, but is subject to considerable uncertainties in the cases of $\bar{p}$ or e^+ . There was an old claim [48] to have observed low-energy $\bar{p}$'s in the cosmic rays: $\phi_{\bar{p}}/\phi_p \sim \text{few} \times 10^{-4}$ for $E < 1$ GeV, but this has now been contradicted by new data that give $\phi_{\bar{p}}/\phi_p < 5.5 \times 10^{-5}$ for $200\text{MeV} < E < 640$ GeV [49]. Some supersymmetric models were proposed [50] to fit the old $\bar{p}$ data: do the new data rule out any realistic sparticle models? Is it realistically possible to get p , e^+ or γ fluxes close to those observed?

We [51] took the local number density n in (25) from a halo model with a local mass density $\rho \sim 0.3 \text{ GeV cm}^{-3}$. The annihilation cross-section σ_{ann} we estimated assuming that the supersymmetric relic particles have $\Omega = 1$, and we estimated the differential fluxes $dn_{\bar{p},e^+,\gamma}/dE$ from each annihilation event using the Lund Monte Carlo which is adapted to fit e^+e^- annihilation data. Typical fluxes [51] from light photinos and higgsinos are shown in Fig. 9. We see that they are considerably below even the new $\bar{p}$ limits. All that one can deduce from these observations are upper limits on the halo density which are about 5 to 10 less than that expected [51], as seen in 10. Also shown in Fig. 10 are the somewhat weaker limits [51] coming from the observed e^+ flux: not shown are the even weaker limits on the halo density from the γ continuum. Since the fluxes are $\propto n^2$ (25), the sensitivity would have to be improved by about two orders of magnitude to have a chance of detecting LSPs in this way, and it is presumably not possible to control the backgrounds down to this level.

It has been suggested [52] that one should look for exclusive relic annihilations to $\gamma\gamma$ final states, which would give a monochromatic line with $E_\gamma \sim m_{\text{relic}}$. Estimates [53] of the possible signal rate, assuming that the LSP is a $\tilde{\gamma}$ with $\Omega h_0^2 \geq \frac{1}{40}$, and that $m_{\tilde{t}} \gg m_{\tilde{\tau}}$, indicate that there is some chance to detect LSPs in this way, but this is not guaranteed.

5.2 HIGH-ENERGY SOLAR NEUTRINOS

These can arise from relic annihilations inside the Sun after relics have been captured. The sequence of events is as follows. A relic particle in a hyperbolic orbit is gravitationally deflected so that it passes through the Sun, where it has some chance of making an elastic collision with a nucleus, and losing sufficient energy to be trapped in an elliptic orbit with its perihelion below the solar surface, so that it makes many collisions on successive passes, and eventually adopts a quasi-isothermal distribution [54] whose effective temperature is fixed by energy balance with the core. (Supersymmetric relics do not cool the core sufficiently to solve the low-energy solar neutrino problem: this ‘‘cosmion’’ solution [55] seems in any case to conflict with indications from the recent SAGE experiment [11] and from helioseismology.) The solar relic population is controlled by evaporation and/or annihilation. In the latter case, some high-energy (typically $E \sim 1$ GeV) neutrinos would be produced that could be detected on Earth [39], either directly in underground proton decay experiments or indirectly via upward-going muons induced by neutrino collisions in lower-lying rock.

The nuclear species in the Sun which dominates the trapping rate is the proton, and so the trapping rate $\Gamma_T \propto \sigma(\chi p \rightarrow \chi p)$ with an exponential dependence on m_χ [56]

$$\Gamma_E \simeq 0.8 \times 10^{-5} x e^{-32.0x} \left(\frac{\sigma(\chi p \rightarrow \chi p)}{10^{-40} \text{ cm}^2} \right) s^{-1} \quad (26)$$

where

$$x \equiv \left(\frac{M_\chi}{3.75 \text{ GeV}} \right) \left(\frac{1.4 \times 10^7 \text{ K}}{T_\chi} \right) \quad (27)$$

There is therefore a critical evaporation mass $m_E : \Gamma_T = \Gamma_E$, with evaporation dominant (negligible) for $m_\chi < (>) m_E$, and $m_E \sim 2\frac{1}{2}$ GeV for a photino. If $m_\chi > m_E$, the annihilation rate Γ_A is given by

$$\Gamma_A = \Gamma_{T/2} \quad (28)$$

and hence is also $\propto \sigma(\chi p \rightarrow \chi p)$.

This results in a high-energy solar neutrino event rate [56,57]

$$R_\nu = 2.7 \times 10^{-2} \theta(m_\chi - m_E) f(m_\chi) \left(\frac{\sigma(\chi p \rightarrow \chi p)}{10^{-40} \text{ cm}^2} \right) \times \left(\frac{\rho_\chi}{0.3 \text{ GeV cm}^{-3}} \right) \left(\frac{300 \text{ kms}^{-1}}{\bar{v}} \right) \left(\sum_{i=1}^3 C_i \right) \text{ events kg}^{-1} \text{ y}^{-1} \quad (29)$$

where $f(m_\chi)$ is a well-known kinematic function, $\bar{v}$ is the mean velocity of the halo particles, and the C_i are known factors derived from $\sigma(\nu_i N)$ and the production rates

for the different neutrino species. Analogously, the high-energy μ event rate is [47, 57]

$$R_\mu = 2.7 \times 10^{-2} \theta(m_\chi - m_E) f(m_\chi) \left(\frac{\sigma(\chi p \rightarrow \chi p)}{10^{-40} \text{ cm}^2} \right) \left(\frac{\rho_\chi}{0.3 \text{ GeV cm}^{-3}} \right) \left(\frac{300 \text{ kms}^{-1}}{\bar{v}} \right) \sum_i C_i' \text{ events kg}^{-1} \text{ y}^{-1} \quad (30)$$

where the C_i' are other known factors.

Before discussing searches for high-energy solar neutrinos and muons, let us first discuss [57] the elastic scattering of relic particles on nucleons, illustrating the analysis by the case of the photino. The photino couples to quarks via virtual squark $\tilde{q}$ exchange, which gives a four-fermion interaction $\tilde{\gamma}\tilde{\gamma}q\bar{q}$ analogous to that in Fermi theory, but scaled by $1/m_{\tilde{q}}^2$. This interaction is the same as that governing annihilation in the early Universe, and the strength, i.e., $m_{\tilde{q}}$, can be fixed by fixing Ω . For elastic scattering, one also needs to know the matrix elements M of the four-fermion operator in a nucleon state. Since the $\tilde{\gamma}$ is a Majorana particle, it couples to spin, so that

$$M \propto \sum_q Q_q^2 \Delta q \quad (31)$$

where the Δq are the fractions of the proton spin carried by different quark species q . The elastic scattering amplitudes for other neutralinos are proportional to other combinations of the Δq [57]. We have some knowledge of these from particle physics measurements, in particular those of β -decay matrix elements in the Fermi theory. Thus neutron decay tells us that [57]

$$\Delta u - \Delta d = 1.26 \quad (32)$$

and hyperon decays combined with flavour $SU(3)$ tell us that [57]

$$\Delta u + \Delta d - 2\Delta s = 0.58 \quad (33)$$

It used to be guessed, on the basis of the naïve quark model (NQM) that $\Delta s = 0$, in which case (32) and (33) suffice to determine Δu and Δd . However, a recent experiment [58] on deep inelastic polarized μp scattering by the EMC indicates that $\Delta s \neq 0$. They

measure

$$\frac{1}{2} \sum_q Q_q^2 \Delta q = \frac{1}{2} \left(\frac{4}{9} \Delta u + \frac{1}{9} \Delta d + \frac{1}{9} \Delta s \right) \simeq 0.126 \pm \quad (34)$$

which can be combined with (32) and (33) to yield [59]

$$\Delta u \simeq 0.77 \quad \Delta d \simeq -0.49 \quad \Delta s \simeq -0.15 \quad (35)$$

In analyzing the high-energy solar ν and μ data, and also in the next section, we will use both the NQM and EMC estimates of the Δq .

Before doing so, however, it is amusing to note that the EMC results (34) imply [59]

$$\sum_q \Delta q = \Delta u + \Delta d + \Delta s \simeq 0.13 \pm 0.24 \quad (36)$$

which is consistent with zero. What then carries the proton spin? A natural explanation is provided by the Skyrme model [60], which predicts that $\sum_q \Delta q = 0$, and that all the proton spin is due to internal angular momentum of its constituents. However, this interpretation is not needed for the analysis of the data on relic particles.

We see from Fig. 11 that, at least if one uses the EMC estimate M for $\sigma(\tilde{\gamma}p \rightarrow \tilde{\gamma}p)$, no region of $m_{\tilde{\gamma}}$ is excluded by the high-energy solar neutrino data, and the best one can deduce is [57]

$$\rho_{\tilde{\gamma}} \lesssim 10 \times \rho_{halo} \text{ for } m_{\tilde{\gamma}} \lesssim 10 \text{ GeV} \quad (37)$$

Similarly, Fig. 11 tells us that [57]

$$\rho_{\tilde{H}} \lesssim 10 \times \rho_{halo} \text{ for } m_{\tilde{H}} \lesssim 20 \text{ GeV} \quad (38)$$

Here there may be an opportunity for MACRO and similar large area detectors looking for upward-going muons.

In the above, we have only discussed high-energy neutrinos from relic annihilations after trapping in the Sun, but one can also consider trapping and annihilation inside the Earth [61]. The Earth loses in the comparison because of its shallower gravitational potential, but gains because it contains heavier nuclei better matched to the mass of the relic, is much closer, and there is a smaller cosmic-ray-induced background. It was found that the Earth even gave larger rates than the Sun if the Higgs boson, which could in principle make a significant contribution to the elastic scattering of relics off

heavy nuclei, were light ($\lesssim 30$) GeV [62]. Unfortunately, this possibility has now been excluded by LEP, which tells us that $m_H > 33$ GeV for the lightest Higgs in the minimal supersymmetric extension of the Standard Model [43]. Therefore the Earth does not at present seem to be a promising place to look for relic annihilations.

5.3 SCATTERING IN THE LABORATORY

Let us first consider elastic scattering on heavy nuclei [63]. The dominant spin-dependent interactions for neutralinos yield an event rate [64]

$$R \simeq 3.9 \left(\frac{100 \text{ GeV}}{m_{\tilde{q}}} \right)^4 M^2 G_{p,n} F_M \left(\frac{\rho}{0.3 \text{ GeV cm}^{-3}} \right) \left(\frac{\Delta \bar{v}}{340 \text{ kms}^{-1}} \right) \text{ kg}^{-1} \text{ d}^{-1} \quad (39)$$

where, as before, we can estimate $m_{\tilde{q}}$ from the requirement that $\Omega = 1$, M is the single-nucleon scattering matrix element discussed in the previous section, for which one may use either the NQM or EMC estimate, $G_{p,n}$ is a factor of 0(1) which corrects for the LSP content when it is not a pure photino [64], F_M is a nuclear figure-of-merit factor to which we return in a moment, ρ is the halo density as before, and $\Delta \bar{v}$ is the relative velocity of the Earth and the relic.

The nuclear figure-of-merit factor F_M may be represented by

$$F_M = Y_{Z,A} \cdot \left(\frac{4m_\chi m_{Z,A}}{(m_\chi + m_{Z,A})^2} \right) \cdot \frac{4}{3} \lambda^2 J(J+1) \quad (40)$$

where $Y_{Z,A}$ is the abundance in the target material of the nuclear species being considered, the parenthesized factor is due to kinematics, we have pulled out an explicit factor of the nuclear spin, and the nuclear physics resides in the constant of proportionality λ^2 .

The first approach that was used [63, 64] to estimate λ^2 for odd-even and even-odd nuclei was based on the naive shell model, according to which the spins of the nucleons are paired, apart from that of the odd nucleon. Then by the Wigner-Eckardt theorem

$$\lambda^2 = \frac{1}{2} \left(1 + \frac{S(S+1) - L(L+1)}{J(J+1)} \right) \quad (41)$$

where $\underline{S}$ and $\underline{L}$ are the nucleon spin and orbital angular momentum respectively. The values of $\underline{S}$ and $\underline{L}$ are shown in the naive shell model, and hence λ^2 can easily be calculated, giving the results [64] shown in Table 1 for some representative nuclei of interest. It is also possible to calculate the nuclear magnetic moments in this model, and

compare them with the experimental values as also shown in Table 1. This comparison serves as a guide to the reliability of the naive shell model, and we see that it works reasonably for model for nuclei where the outermost shell is almost complete (one hole) or only contains one nucleon.

The nuclear figures of merit F_M (40) calculated in the naive shell model [63, 59] are shown in Fig. 12. Using these, we found [64] the event rates shown in Table 2: they are typically between 1 and 10^{-2} events $\text{kg}^{-1} \text{ day}^{-1}$.

Since these first naive shell model calculations, various strategies have been used to calculate improved estimates. One of these is to calculate [65] corrections to the above naive single-particle estimates using a complete shell model basis of states. This becomes unwieldy for heavier nuclei, and results have been given only for ${}^7\text{Li}$, ${}^9\text{Be}$, ${}^{11}\text{Be}$, ${}^{17}\text{O}$, ${}^{19}\text{F}$ and ${}^{35}\text{Cl}$. One favoured simplifying hypothesis is that all the nuclear spin is carried collectively by the species of nucleon whose number is odd, with λ^2 determined by the nuclear magnetic moment whose number is odd: the odd-group model. This idea has been applied [66] to ${}^7\text{Li}$, ${}^{11}\text{B}$, ${}^{19}\text{F}$, ${}^{27}\text{Al}$, ${}^{35}\text{Cl}$, ${}^{39}\text{K}$, ${}^3\text{He}$, ${}^{13}\text{C}$, ${}^{17}\text{O}$, ${}^{21}\text{Ne}$, ${}^{25}\text{Mg}$ and ${}^{29}\text{Si}$.

The shell-model predictions [59] for F_M shown in Fig. 12 are in general larger than the odd-group model predictions, in particular for heavier nuclei. However, we do note that the odd-group and naive shell-model predictions for ${}^{19}\text{F}$ almost coincide. Indeed, all the above models give similar results for this nucleus: the even and odd angular momenta in these different models are:

	naive[64]	complete[65]	odd - group[66]	mirror[66]
S_{even}	0	0.109	0	-0.047
S_{odd}	0.5	0.441	0.38	0.415

Hence we regard the rate calculations for ${}^{19}\text{F}$ as being particularly reliable.

To make improved rate calculations for ${}^{19}\text{F}$, we [59] (1) choose $m_{\tilde{q}} : \Omega = 1$ for $h_0 = 1/2$, (2) use the latest EMC spin-dependent matrix elements for protons and neutrons, (3) use the odd-group model for the nuclear figure-of-merit, (4) incorporate correctly gaugino/higgsino mixing, (5) take into account LEP constraints [46] on the LSP parameters, (6) restrict our attention to $m_\chi < m_W$. We find rates above 10^{-2} events $\text{kg}^{-1} \text{ day}^{-1}$ in much of the LSP parameter space.

Present searches for dark matter particles have been using ionization detectors such as ${}^{76}\text{Ge}$ [18]. Unfortunately, the materials used mainly contain nuclei with no spin, and hence are not good for LSP searches. Other ideas for future detectors involve detecting the nuclear recoil energy of $0(\text{keV})$ by techniques such as direct phonon detection,

calorimetry at $\sim 100m_K$, or a normal/superconducting transition in a resistance, tunnel junction, or granules [2]. The optimal future search strategy is likely to combine ionization with one of these other techniques.

In order to improve the event signature, some consideration has been given to inelastic scattering $\tilde{\gamma}N \rightarrow \tilde{\gamma}N^*$ followed by $N^* \rightarrow N\gamma$ decay [67]. Thus we have a double signature of nuclear recoil energy and a (10 to 50) keV de-excitation photon. One could perhaps even hope to do such a double experiment at room temperature. Unfortunately, the inelastic scattering rates calculated [67] in the naive shell model are discouragingly small, as seen in Table 3. The reasons for this are that the excitation energies are high, entailing a kinematical suppression, and that the transition magnetic moments are small. Since typical rates are $O(10^{-4})$ events $\text{kg}^{-1}\text{day}^{-1}$, one would need a detector weighing ~ 1 ton. Such a large detector could presumably not be cryogenic, and one therefore might need to master room-temperature techniques.

6. SHADOW MATTER

This is a conjectured new set of particles that have their own “hidden” gauge interactions among themselves, but have in general only Planck-strength $(1/m_P^2)$ interactions with ordinary matter [68]*. Shadow matter is present in many low-energy models obtainable from string. For example, in the old Calabi-Yau compactifications [69] the gauge group comprised an observable part contained in E_6 and a hidden gauge group contained in E_8 . More recently, a flipped $SU(5)$ model derived [70] from a string theory with free fermions on the world-sheet has a priori an $[SU(5) \times U(1)] \times U(1)^4 \times [SU(4) \times SO(10)]$ gauge group, where the two sets of square parentheses contain the observable and hidden gauge interactions, and the other $U(1)^4$ factors connect observable and shadow matter.

If the hidden gauge sector only contains gauge supermultiplets, all the hidden sector particles are expected to be unstable, with hidden-gluonball decays to $\gamma\gamma$ or $g\bar{g}$, and hidden-gluonballino decays to $\tilde{\gamma}\tilde{\gamma}, g\bar{g}$, etc. [71]. This was the situation in the old Calabi-Yau models with an unbroken E_8 hidden group [69].

However, the presence of chiral matter supermultiplets in the shadow sector now seems to be generic. For example, in the flipped $SU(5)$ model mentioned above there are $\underline{4}$ and $\underline{6}$ representations of hidden $SU(4)$, and $\underline{10}$ representations of hidden $SO(10)$, with masses $m \ll m_P$ [70]. It is even possible that such hidden sector matter representations may play a key role in breaking supersymmetry and communicating this breaking to

* Although, as we shall see, some new variants of shadow matter can have electric charges.

the observable sector via the connecting $U(1)$'s [72]. They may also be the only way to avoid observable fractionally-charged states [70]. All existing low-energy string-derived models are based on world-sheet currents with a level-one Kac-Moody algebra. In this case, it has been shown that the light ($m \ll m_P$) spectrum must contain fractional charges [73], e.g., $Q_{em} = \frac{1}{n}$ in Z_n orbifold models. The only way to avoid seeing them is to confine them via hidden-sector gauge interactions [70,4], by analogy with the confinement of fractionally charged quarks in QCD. We call [4] the resulting integer-charged shadow matter particles “cryptons”, which may include two-constituent mesons, three-constituent baryons, four-constituent “tetrons”, etc.

Let us consider in more detail as an example the flipped $SU(5)$ model [70] mentioned above. Its hidden matter representations transform under $[SU(4) \times SO(10)]_{Q_{em}}$ as

$$\bar{F} \equiv (\underline{\bar{3}}, 1)_{\pm \frac{1}{2}, 0}; \quad \Delta \equiv (\underline{6}, 1)_0, \quad T \equiv (1, \underline{10})_0 \quad (43)$$

Confinement then yields the following cryptons [4]:

$$\left. \begin{array}{ll} \text{mesons:} & (TT)_0, \quad (\Delta\Delta)_0, \quad (F\bar{F})_{0,\pm 1} \\ \text{baryons:} & (FFF\Delta)_{0,\pm 1}, \quad (\bar{F}\bar{F}\bar{F}\Delta)_{0,\pm 1} \\ \text{tetrons:} & (FFFF)_{0,\pm 1,\pm 2}, \quad (\bar{F}\bar{F}\bar{F}\bar{F})_{0,\pm 1,\pm 2} \end{array} \right\} \quad (44)$$

We expect the baryons and tetrons to have masses $m \sim \Lambda_{SU(4)} \sim 10^{12}$ GeV, whereas some mesons could be much lighter because of chiral symmetries. In this particular model the lightest tetron would be electrically neutral [4], but it would in general be charged. We would expect at least some of the states (44) to be metastable, decaying only via non-renormalizable interactions, analogously to the proton in GUTs. Therefore there could be one or more charged or neutral species of relic shadow matter particles left over from the Big Bang.

Their relic abundance is very uncertain, and has no natural reason to give $\Omega \sim 1$. In the absence of any substantial entropy production after crypton formation, one would expect [74,4]

$$\Omega_{\text{cryptons}} \sim \left(\frac{\Lambda_{SU(4)}}{100 \text{ TeV}} \right)^2 \quad (45)$$

which would be unacceptable for $\Lambda_{SU(4)} \sim 10^{12}$ GeV as we expect. However, there are plenty of plausible sources of entropy production, including inflation, gravitino decay with $\tau_{3/2} \sim m_{\tilde{P}}^2/m_{3/2}^3$, decays of coherent scalar fields along a generic “flat” direction of the effective potential, and the decays of other long- (but shorter-) lived

hidden sector particles [4]. Neutral shadow matter particles would be very difficult to detect or exclude. However, charged shadow matter particles may well already be experimentally excluded for masses $\lesssim 10^{16}$ GeV by the survival of neutron stars [75]. If not, they should be detectable if $\Omega \sim 1$, and could even constitute a "smoking gun" for string.

7. A CHALLENGE

I have discussed many candidates for dark matter. The challenge to experimental physicists is to detect one of them, and thereby complete the Copernican Revolution. We learnt several centuries ago that we are not the centre of the Universe. Now astrophysicists tell us that we are not made of the same stuff as most of the matter in the Universe. The task of particle physicists is to prove it!

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FIGURE CAPTIONS

Fig. 1 : Values of mass and Z^0 coupling strength $\sin^2\theta_Z$ relative to a conventional neutrino consistent with LEP [19] and giving $\Omega = 1$ for $h_0 = 0.5$ or 0.7 (a) for Majorana fermions and (b) for Dirac Fermions. In the latter case the constraint [18] from underground ^{76}Ge experiments is also shown.

Fig. 2 : The relative values of $\hat{f}_a$, normalized so that $\hat{f}_a = 1$ for $\beta = 0$ for the NQM model ($\hat{f}_a = f_a/0.19 \times 10^{12} \text{ GeV}$), obtained [29] requiring equal energy loss to axions and neutrinos, plotted versus the angle $\beta : \tan\beta = v_2/v_1$ for NQM and EMC models.

Fig. 3 : The lower bound [31] on the parameter f_Φ is displayed as a function of the chemical potential μ reached in the deconfined core of supernova 1987A. This astrophysical constraint, derived from the requirement that axion luminosity from transverse gluon decay should not exceed $2 \times 10^{53} \text{ erg}$ in the supernova event, is stronger for scalar (solid) than for pseudoscalar (short dash) bosons.

Fig. 4 : Search for halo axions [33].

Fig. 5 : Limits on the abundances of heavy isotopes of various elements. The curves are labeled by the chemical element studied [38]. The dashed lines are preliminary results.

Fig. 6 : Regions of the (μ, M_2) plane excluded [46] by LEP, CDF and UA2 data, for the case $\tan\beta = 2$.

Fig. 7 : Domains of the (μ, M_2) plane consistent [19] with LEP, CDF and $\frac{1}{4} < \Omega h_0^2 < \frac{1}{2}$ for some choice $120 \text{ GeV} \leq m_{\tilde{q}} \leq 1 \text{ TeV}$, $39 \text{ GeV} \leq m_A \leq 1 \text{ TeV}$ and $1.5 < \tan\beta < 10$.

Note the favoured regions (1), (2), (3) and (4) discussed in the text and the rays (a), (b), (c) and (d).

Fig. 8 : The LSP mass density as a function of m_{LSP} along the ray (a) $\mu/M_2 = -5$ shown in Fig. 7 [19].

Fig. 9 : Ratios of $\tilde{p}$ and p fluxes from different relic candidates [51]. Also shown is a recent upper limit on the $\tilde{p}$ flux. Previous measurements at higher energies are off-scale beyond the top of the figure. The curves are calculated with the default options $\rho_{\tilde{\chi}}^h = 0.3 \text{ GeV cm}^{-3}$ and $\tau_{\tilde{p}} = 2 \times 10^{15} \text{ s}$. The solid/dotted/dashed/dash/dotted curves are for $m_{\tilde{q}} = 3 \text{ GeV}/m_{\tilde{q}} = 6 \text{ GeV}/m_{\tilde{q}} = 10 \text{ GeV}/m_{\tilde{q}} = 15 \text{ GeV}$.

Fig. 10 : Upper limits on $\rho_{\tilde{\chi}}^h \sqrt{\tilde{p}}$ from the $\tilde{p}$ flux [51], for $\chi = \tilde{\gamma}$ and $\tilde{h}$ as functions of $m_{\tilde{\chi}}$. Also shown are "upper limits" on $\rho_{h_{\tilde{\chi}}} \sqrt{\tilde{e}^+}$ from measurements of the e^+ cosmic ray flux above 2 GeV . Lower-energy e^+ data are not competitive for constraining relics in the mass ranges considered here.

Fig. 11 : Upper limits on $\rho_{\tilde{\chi}}$ [57]: (a) $\tilde{\chi} = \tilde{\gamma}$, using Fréjus data under different model assumptions (A), (B) and (C), and IMB data under assumption (A); (b) $\tilde{\chi} = \tilde{h}$ using Fréjus data under model assumptions (E), (F). The $\tilde{E}$ and $\tilde{F}$ curves are obtained by removing from the rate a model-dependent factor of $\cos^2(2\beta)$. The horizontal lines indicate the astrophysically preferred value and probable limits on $\rho_{\tilde{\chi}}$.

Fig. 12 : Figures of merit calculated [59] for different nuclei using the shell model.

Table 1

Typical rates in events/kgday for different elements, calculated using the shell-model and NQM (top lines) or EMC (bottom lines) nucleon scattering matrix elements.

Isotope	m_{LSP}	4	8	15	20
1_1H		0.29	0.0085	0.0013	0.00046
2_1H		0.89	0.026	0.0037	0.0013
3_1H		0.002	0.0029	0.00050	0.00019
4_2He		2.8	0.33	0.051	0.025
6_3Li		1.4	0.18	0.032	0.012
7_3Li		0.017	0.0030	0.00071	0.00033
${}^{10}_5B$		1.0	0.18	0.029	0.016
${}^{11}_5B$		3.1	0.55	0.12	0.048
${}^{12}_6C$		0.66	0.14	0.030	0.013
${}^{13}_6C$		0.0086	0.0021	0.00068	0.00035
${}^{16}_8O$		0.016	0.0035	0.00090	0.00039
${}^{18}_8O$		2.0	0.40	0.92	0.052
${}^{20}_{10}Ne$		0.84	0.18	0.0044	0.019
${}^{22}_{10}Ne$		2.6	0.54	0.13	0.055
${}^{24}_{12}Mg$		1.2	0.29	0.083	0.038
${}^{26}_{12}Mg$		3.5	0.87	0.25	0.12
${}^{28}_{14}Si$		0.42	0.11	0.036	0.018
${}^{30}_{14}Si$		1.3	0.33	0.11	0.051
${}^{32}_{16}S$		0.23	0.068	0.026	0.014
${}^{34}_{16}S$		0.71	0.21	0.076	0.040
${}^{68}_{31}Ga + {}^{71}_{31}Ga$		0.23	0.070	0.028	0.015
${}^{70}_{31}Ga + {}^{73}_{31}Ga$		0.70	0.21	0.082	0.045
${}^{75}_{35}As$		0.21	0.066	0.027	0.015
${}^{77}_{35}Br + {}^{81}_{35}Br$		0.66	0.20	0.079	0.043
${}^{79}_{35}Br + {}^{83}_{35}Br$		0.20	0.063	0.026	0.014
${}^{203}_{81}Tl + {}^{205}_{81}Tl$		0.62	0.19	0.075	0.041
${}^{207}_{81}Tl + {}^{209}_{81}Tl$		0.15	0.050	0.022	0.013
${}^{209}_{81}Tl$		0.46	0.15	0.066	0.038

Table 2
Shell Model Factors

The odd group model may give more reliable results for λ^2 in cases where $\mu_{Th} \neq \mu_{Ez}$.

A Z Isotope	J	Abundance	Shell Model	μ_{Th}	μ_{Ez}	$\lambda^2 J(J+1)$
1_1H	1/2	100*	$S_{1/2}$ proton	2.79	2.79	3/4
7_3Li	3/2	100 ^c	$P_{3/2}$ proton	3.79	3.256	5/12
${}^{11}_5B$	3/2	100 ^c	$P_{3/2}$ proton hole	3.79	2.689	5/12
${}^{15}_7N$	1/2	0.37	$P_{1/2}$ proton hole	-0.26	-0.283	1/12
${}^{19}_9F$	1/2	100	$S_{1/2}$ proton	2.79	2.629	3/4
${}^{27}_{13}Al$	5/2	100	$D_{5/2}$ proton hole	4.79	3.642	7/20
${}^{35}_{17}Cl$	3/2	100 ^c	$D_{3/2}$ proton	0.13	0.822	3/20
${}^{51}_{23}V$	7/2	99.8	$F_{7/2}$ proton	5.79	5.151	9/28
${}^{69}_{31}Ga, {}^{71}_{31}Ga$	3/2	60.1 39.9	$P_{3/2}$ proton hole	3.79	2.917 2.562	5/12
${}^{75}_{33}As$	3/2	100	$P_{3/2}$ proton hole	3.79	1.439	5/12
${}^{79}_{35}Br, {}^{81}_{35}Br$	3/2	50.7 49.3	$P_{3/2}$ proton hole	3.79	2.106 2.271	5/12
${}^{93}_{41}Nb$	9/2	100	$G_{9/2}$ proton	6.79	6.171	11/36
${}^{107}_{47}Ag, {}^{109}_{47}Ag$	1/2	57.8 48.2	$P_{1/2}$ proton	-0.26	-0.114 -0.151	1/12
${}^{121}_{51}Sb, {}^{123}_{51}Sb$	5/2	57.3 42.7	$G_{7/2}$ proton	1.71	3.359 2.547	7/36
${}^{129}_{53}I$	5/2	100	$D_{5/2}$ proton	5.79	2.808	5/20
${}^{133}_{55}Cs$	7/2	100	$G_{7/2}$ proton	1.71	2.579	7/36
${}^{139}_{57}La$	7/2	99.9	$G_{7/2}$ proton hole	1.71	0.50	7/36
${}^{203}_{81}Tl, {}^{205}_{81}Tl$	1/2	29.5 70.5	$S_{1/2}$ proton	2.79	1.622 1.338	3/4

FIGURE 1

FIGURE 2

FIGURE 3

FIGURE 4

FIGURE 5

FIGURE 6

FIGURE 7

FIGURE 8

FIGURE 9

FIGURE 10

FIGURE 11

FIGURE 12