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IMPLICATIONS FOR DARK MATTER PARTICLES OF SEARCHES FOR ENERGETIC SOLAR NEUTRINOS

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ABSTRACT

We investigate whether recent searches for energetic solar neutrinos give new constraints on candidate dark matter particles such as photinos $\tilde{\gamma}$ and higgsinos $\tilde{h}$, in view of the astrophysical and particle physics uncertainties. Previous quark model estimates of the elastic cross-sections responsible for $\tilde{\gamma}$ or $\tilde{h}$ capture in the Sun can be improved using hyperon decay data, but are thrown into doubt by recent European Muon Collaboration data on the spin-dependent proton structure function, which reduce the estimate of $\sigma(\tilde{\gamma}p)$ by a factor up to four. We emphasize the need to include charm and bottom quark fragmentation effects when calculating the flux of neutrinos from dark matter annihilation events, which may reduce the observable flux by a factor ~ 2 . Even the most stringent upper limits from the Fréjus experiment still allow a local $\tilde{\gamma}$ or $\tilde{h}$ density considerably above the range of 0.20 to 0.43 GeV/cm³ favoured by astrophysics, and no range of $m_{\tilde{\gamma}}$ or $m_{\tilde{h}}$ can yet be excluded.

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ERRATUM

First, the numerical prefactor in Eq. (1) was written as 1.8×10^{-2} , but should be 2.7×10^{-2} . Second, in Eq. (23) there is a missing factor of $m_{\tilde{\nu}}$, which is most conveniently placed directly after the factor, $\sigma(\tilde{\gamma}p \rightarrow \tilde{\gamma}p)$, and before the sums. Equations (23) should therefore read:

$$\begin{aligned}
 R_{\mu} \simeq & 1.1 \times 10^{-2} \left(\frac{\text{Area}}{100 \text{ m}^2} \right) \left(\frac{\rho_L}{0.3 \text{ GeV cm}^{-3}} \right) \left(\frac{300 \text{ km s}^{-1}}{\bar{v}} \right) \\
 & \times \Theta(m_{\tilde{g}} - m_E) f(m_{\tilde{g}}/m_c) \left(\frac{\sigma(\tilde{g}p \rightarrow \tilde{g}p)}{10^{-40} \text{ cm}^2} \right) m_{\tilde{g}} \\
 & \times \sum_i a_i b_i \sum_f B_f \langle N z^2 \rangle_{fi} : b_{\nu_{\mu}} = 0.5, b_{\bar{\nu}_{\mu}} = 0.67
 \end{aligned} \tag{23}$$

Since these are purely typographical errors, all the results of the paper are unaffected.

There is an emerging consensus that non-dissipative dark matter is prevalent in galactic halos¹⁾. One of the prime possibilities is that the halo dark matter consists of some species of massive, non-relativistic, neutral, weakly-interacting particle. A possible candidate for this dark matter particle is the lightest supersymmetric particle (LSP), which could be a photino $\tilde{\gamma}$, higgsino $\tilde{h}$ or sneutrino $\tilde{\nu}$ ²⁾. Methods of detecting dark matter particles are now being actively pursued³⁾. The most convincing would be direct detection of the elastic scattering of dark matter particles off nuclei in the laboratory⁴⁾, unsuccessful searches for which have already been used to exclude sneutrinos in the range $16 \text{ GeV} < m_{\tilde{\nu}} < 1 \text{ TeV}$ ⁵⁾. An elegant but less direct method⁶⁾ is based on the observation that certain candidate dark matter particles have some probability of being trapped in the Sun after losing energy through elastic scattering^{7),8)}. Evaporation is negligible for masses above $2\frac{1}{2} \text{ GeV}$ ^{9),10)}, in which case the solar population is stabilized by annihilations. Such annihilations, e.g., $\tilde{\gamma}\tilde{\gamma} \rightarrow f\bar{f}$, can produce energetic neutrinos via the prompt decays of heavy flavours: $f = \tau, c, b, t$. This idea has been pursued extensively¹¹⁾⁻¹⁵⁾ and varying estimates given of the resulting ν flux.

Data are now becoming available from underground detectors¹⁶⁾⁻¹⁸⁾, most recently the Fréjus experiment¹⁸⁾, both on contained ν interactions and on upward-going muons induced by ν_{μ} interactions in rock. These provide upper limits on the energetic solar ν flux which can be confronted with different dark matter predictions. Some candidates seem to be clearly excluded by such analyses, e.g., $\tilde{\nu}_{\mu}$ and $\tilde{\nu}_e$ with masses $\gtrsim 4 \text{ GeV}$ ^{13),14)}. However, the confrontation is much more delicate for some candidates, such as the $\tilde{\gamma}$ and $\tilde{h}$, for which the experimental upper limits can be comparable with the ν event estimate. For $\tilde{\gamma}\tilde{\gamma}$ annihilation, this is approximately

$$R_{\nu} \approx 1.8 \times 10^{-2} \left(\frac{\rho_L}{0.3 \text{ GeV cm}^{-3}} \right) \left(\frac{300 \text{ km s}^{-1}}{\bar{v}} \right) \theta(m_{\tilde{\gamma}} - m_E) f(m_{\tilde{\gamma}}/m_c) \times$$

$$\times \left(\frac{\sigma(\tilde{\gamma}p \rightarrow \tilde{\gamma}p)}{10^{-40} \text{ cm}^2} \right) \times \sum_i a_i \sum_f B_f \langle N_z \rangle_{fi} \quad \text{events/ktony} \quad (1)$$

for contained ν events, where ρ_L is the local halo density, $\bar{v}$ is the three-dimensional velocity dispersion of the dark matter particles, m_c is a critical mass⁸⁾, m_E is an evaporation mass $\sim 2\frac{1}{2} \text{ GeV}$ ¹⁰⁾, the a_i are simple numerical factors in the cross-section for ν_i interactions which depend on the composition of the detector, B_f is the branching ratio for $\tilde{\gamma}\tilde{\gamma} \rightarrow f\bar{f}$ annihilation, N is the number of neutrinos

produced in each annihilation and z the scaled energy fraction $E_{\nu_i}/m_{\tilde{\gamma}}$. In this paper we refine previous numerical estimates^{6),11)-15)} based on (1), survey the inherent uncertainties, and study whether the present data¹⁶⁾⁻¹⁸⁾ can be used to give significant constraints on $m_{\tilde{\gamma}}$ or $m_{\tilde{h}}$.

Among the uncertainties we discuss are the following. (a) The principal uncertainty in the numerical prefactor in (1) is due to the shape of the dark matter velocity distribution, but is not very large. (b) Model studies¹⁹⁾ indicate that the local dark matter density ρ_L is probably between 0.2 and 0.43 GeV/cm³, with a preferred value $\rho_L = 0.3$ GeV/cm³. (c) The best estimate of $\bar{v}$ is 300 kms with an error of about 10%. (d) The factor written as $f(m_c/m_{\tilde{\gamma}})$ is often approximated⁷⁾ by $\min(1, m_c/m_{\tilde{\gamma}})$, but is better approximated by a calculation due to Gould⁸⁾. (e) Estimates of $\sigma(\tilde{\gamma}p \rightarrow \tilde{\gamma}p)$ depend principally on estimates of the squark masses and on the proton matrix elements of $\bar{q}\gamma_{\mu}\gamma_5 q$ operators. The overall scale of squark and slepton masses $m_{\tilde{q}}$ and $m_{\tilde{\chi}}$ can be fixed by requiring the relic photino mass density $\rho_{\tilde{\gamma}}$ to equal the critical density ρ_c required for closure: $\Omega_{\tilde{\gamma}} \equiv \rho_{\tilde{\gamma}}/\rho_c = 1$ ²⁰⁾. However, the ratios $m_{\tilde{q}}/m_{\tilde{\chi}}$ are in general model-dependent and $\neq 1$. Previous estimates^{4),6),11)-14),21)} of $\langle p|\bar{q}\gamma_{\mu}\gamma_5 q|p\rangle$ have mainly been based on the naïve quark model, but these can be refined using experimental data. The matrix elements of flavour-octet currents can be estimated on the basis of hyperon β -decay measurements²²⁾, which differ somewhat from the naïve quark model estimates. However, until recently the matrix element of the singlet current had not been measured, and it was generally fixed by assuming that $\langle p|\bar{s}\gamma_{\mu}\gamma_5 s|p\rangle \approx 0$ ^{23),24)}. This assumption has recently been called into question by European Muon Collaboration (EMC) data²⁵⁾ on the spin-dependent proton structure function $g_1^p(x, Q^2)$, whose integral is in theory proportional to the same combination of $\langle p|\bar{q}\gamma_{\mu}\gamma_5 q|p\rangle$ as the $\tilde{\gamma}p$ scattering matrix element. The EMC value of the integral at $Q^2 \gtrsim 3$ GeV² is about $\frac{1}{2}$ the naïve quark model prediction, and indicates that $\langle p|\bar{s}\gamma_{\mu}\gamma_5 s|p\rangle \neq 0$. The $\tilde{\gamma}p$ scattering process explores these matrix elements at scales $\sim m_{\tilde{q}}^2 \sim 10^4$ GeV², but we do not expect significant variation between $Q^2 \sim 3$ GeV² and these scales, though there might be significant Q^2 -dependence at lower Q^2 ²⁶⁾. The EMC data suggest the possibility that previous estimates^{4),6),11)-14),21)} of $\sigma(\tilde{\gamma}p \rightarrow \tilde{\gamma}p)$ may need to be reduced by a factor ~ 4 . Finally, we emphasize¹⁵⁾ that quark fragmentation effects should be taken into account in estimating $\langle Nz \rangle_{fi}$ in (1), and that these reduce naïve estimates^{6),11)-14)} of c and b quark contributions by factors of 2 to 3. For completeness, we also present corresponding results for higgsinos $\tilde{h}$.

We use the recent Fréjus data¹⁸⁾ to give upper limits on the local photino density $\rho_{\tilde{\gamma}}$ under different theoretical assumptions on $m_{\tilde{\chi}}/m_{\tilde{q}}$, on Ω and on the

$\langle p | \bar{q} \gamma_\mu \gamma_5 q | p \rangle$ matrix elements. These are generally above the astrophysical estimate¹⁹⁾ of ρ_L except under extreme assumptions, so that there is no safe bound on $m_{\tilde{\gamma}}$ or $m_{\tilde{h}}$ yet. However, future data on upward-going muons may soon sharpen considerably the confrontation between theory and experiment.

The rate for trapping of cold dark matter particles $\tilde{X}$ by the Sun is^{7),8)}

$$\Gamma_T = a \left(\frac{3\pi}{2} \right)^{\frac{1}{2}} N_\odot \frac{v_\odot^2}{\bar{v}} n_{\tilde{X}} f_G \sum_i x_i \sigma(\tilde{X}i \rightarrow \tilde{X}i) f(m_{\tilde{X}}/m_{c_i}) \quad (2)$$

where $N_\odot$ is the number of nuclei in the Sun, $v_\odot = 617 \text{ kms}^{-1}$ is the escape velocity from the surface of the Sun, $\bar{v}$ is the local velocity dispersion in the halo, $n_{\tilde{X}}$ is the number density of $\tilde{X}$ particles in the solar neighbourhood, and x_i is the fractional solar abundance by number of nuclear species i . The factor $a = 0.89$ for a nuclear species whose number density traces the mass density of the Sun⁷⁾. Since $\tilde{\gamma}$ and $\tilde{h}$ do not have coherent cross-sections for scattering off heavy nuclei, the sum over species i in (2) can be truncated to protons p , for which this tracing assumption is good and $x_p = 0.86$. The factor $f_G \approx 1.5$ corrects⁸⁾ the estimate of Ref. 7) for the motion of the Sun through the galactic halo and for a numerical slip. The remaining kinematic factors $f(m_{\tilde{X}}/m_{c_i})$ express the mismatching of the $\tilde{X}$ and nuclear masses. These are often approximated⁷⁾ by

$$f(m_{\tilde{X}}/m_{c_i}) = \min(1, m_{\tilde{X}}/m_{c_i}) \quad (3a)$$

where

$$m_{c_i} = \frac{1}{3\beta} \left(\frac{m_i^2 + m_{\tilde{X}}^2}{m_i} \right) \frac{\bar{v}^2}{v_\odot^2} \quad : \quad \beta \approx 3.4 \quad (3b)$$

but this is not⁸⁾ very accurate for $m_{\tilde{X}}/m_{c_i} \sim 1$. In the case of protons, a better approximation¹⁵⁾ is

$$f(m_{\tilde{X}}/m_c) = \left(\frac{m_c^a}{m_c^a + m_{\tilde{X}}^a} \right)^{1/a} : a=1.5, m_c = 20.4 \text{ GeV} \quad (4)$$

which is accurate to within ~5%.

The population of relic particles $\tilde{X}$ in the Sun is controlled by evaporation⁹⁾ and annihilation. The evaporation rate is calculated¹⁰⁾ to be

$$C_E = 0.8 \times 10^{-5} x^2 e^{-32.9x} \left(\frac{\sigma(\tilde{X}p \rightarrow \tilde{X}p)}{10^{-40} \text{ cm}^2} \right) \text{ s}^{-1} \quad (5a)$$

where

$$\chi = \left(\frac{m_{\tilde{X}}}{3.75 \text{ GeV}} \right) \left(\frac{1.4 \times 10^7 \text{ }^\circ\text{K}}{T_{\tilde{X}}} \right) \quad (5b)$$

with $T_{\tilde{X}}$ the effective temperature of the isothermal distribution of $\tilde{X}$ particles within the Sun. Because of the rapidly varying exponential factor in (5a), evaporation dominates over annihilation for $m_{\tilde{X}}$ below some critical evaporation mass m_E , but can be neglected by comparison to annihilation when $m_{\tilde{X}} > m_E^*$). This explains the $\theta(m_{\tilde{X}} - m_E)$ factor in (1), where the evaporation mass $m_E \sim 2\frac{1}{2}$ GeV for $\tilde{X} = \tilde{\gamma}$ and $\tilde{h}$ using our estimates of $\sigma(\tilde{X}p \rightarrow \tilde{X}p)$ and $T_{\tilde{X}}$. When $m_{\tilde{X}} > m_E$, the annihilation rate can be deduced from $\Gamma_T(2)$ by dividing by two for self-conjugate particles such as the $\tilde{\gamma}$ and $\tilde{h}$, and the resulting flux at the Earth is obtained by dividing by $4\pi R^2$, where R is the Earth-Sun distance. In this way, we obtain the solar contributions to the prefactors in the observable neutrino interaction rate (1).

The main astrophysical uncertainty in (1) is in ρ_L . If one assumes¹⁹⁾ that galaxies form in pre-existing potential wells of dark matter, then an adiabatic model²⁷⁾ can be used to bound ρ_L . The main constraints are that the circular velocity in the Galaxy not fall below 200 kms⁻¹ at large radii ~ 60 kpc, and that the ratio F of luminous to total mass inside the region where the gas cools to form visible galaxies obey $F \gtrsim 0.05$. One then obtains¹⁹⁾

$$0.20 \lesssim \rho_L / \text{GeVcm}^{-3} \lesssim 0.43 \quad (6a)$$

if the shape of the potential well before galaxy formation is that expected for cold dark matter^{**)}. For more likely values of the parameters

$$0.30 \lesssim \rho_L / \text{GeVcm}^{-3} \lesssim 0.43 \quad (6b)$$

and the typical estimate $F \approx 0.1$ implies $\rho_L \approx 0.3 \text{ GeVcm}^{-3}$. Needless to say, these are dynamical bounds and the lower bound therefore applies to particle dark matter only if it makes up all the mass of the halo. If some of the halo dark matter is, e.g. little Jupiters or black holes²⁸⁾, then the particle dark matter density could be below the lower bounds (6). Also, meaningful confidence levels cannot be assigned to the bounds on ρ_L because the systematic uncertainties are large. For

*) The width of the transition region is ~ 200 MeV, and is ignored in our subsequent discussion.

**) The lower bound is relaxed by only 15% if one relaxes this assumption.

these reasons, we will prefer to use the experimental constraints to provide upper bounds on the particle dark matter contribution to ρ_L .

Next comes the uncertainty in $\bar{v}$. If one assumes that the velocity dispersion of the halo particles is homogeneous and isotropic, i.e., the halo is an isothermal sphere, then "hydrostatic" equilibrium implies $\bar{v}^2 = (3/2)v_\infty^2$, where v_∞ is the asymptotic circular velocity. If one assumes a constant circular velocity outside the solar radius, then $v_\infty = v_c$, the local circular velocity for which the IAU recommended value²⁹⁾ is $v_c = 220 \text{ kms}^{-1}$ and a straight mean of the observations yields²⁹⁾ $v_c = 222 \pm 20 \text{ kms}^{-1}$, implying $\bar{v} = 270 \pm 25 \text{ kms}^{-1}$. If one does not assume a constant circular velocity, but assumes that $v_\infty = v(60 \text{ kpc})$, observations suggest that probably $200 \lesssim v_\infty/\text{kms}^{-1} \lesssim 260^*)$, implying $245 \lesssim \bar{v}/\text{kms}^{-1} \lesssim 320$. The halo particles cannot have an exactly isothermal distribution, because they are compressed during galaxy formation²⁷⁾ and distorted by the potential of the disc. These effects on $\bar{v}$ have yet to be quantified, but appear to be small³²⁾. Motivated by the fact that higher $\bar{v}$ implies lower $R_\nu(1)$, we will take $\bar{v} = 300 \text{ kms}^{-1}$ for our analysis. It is justified to take the highest value of $\bar{v}$ that is allowed by the uncertainty in v_c in deriving a conservative bound on masses.

Next we turn to uncertainties in estimating $\sigma(\tilde{\gamma}p \rightarrow \tilde{\gamma}p)$, which takes the following form^{6),21)} if the supersymmetric partners of the left- and right-handed squarks have equal masses, $m_{\tilde{q}_L} = m_{\tilde{q}_R} = m_{\tilde{q}}$:

$$\sigma(\tilde{\gamma}p \rightarrow \tilde{\gamma}p) = \frac{3}{\pi} \frac{m_{\tilde{\gamma}}^2 m_p^2 (4\pi\alpha)^2}{(m_{\tilde{\gamma}} + m_p)^2} M^2 : M = \sum_q \frac{Q_q^2}{m_q^2} \Delta q \quad (7a)$$

where $Q_q = 2/3$ for u, c , $-1/3$ for d, s and Δq is defined by

$$\langle p | \bar{q} \gamma_\mu \gamma_5 q | p \rangle = \Delta q \cdot 2s_\mu \quad (7b)$$

where s_μ is the proton spin. Differences between $m_{\tilde{q}_L}$ and $m_{\tilde{q}_R}$ can arise from different renormalizations of these parameters by interactions at scales $\gtrsim m_W$, and from differences in Higgs vacuum expectation values, but we disregard such possibilities here. The overall scale of the sparticle masses as functions of $m_{\tilde{\gamma}}$ can be fixed by

*) The Magellanic stream³⁰⁾ indicates a high $v(60 \text{ kpc}) \approx 240 \text{ kms}^{-1}$, whilst globular cluster data³¹⁾ indicate $\approx 200 \text{ kms}^{-1}$.

requiring closure density for photinos: $\Omega_{\tilde{\gamma}} = 1^{*})$. If $m_{\tilde{\chi}} = m_{\tilde{q}}$ (EM), one needs masses^{6),13),14)}:

$$m_{\tilde{q}} \simeq 56 \text{ GeV} (1 + 0.04 m_{\tilde{g}}^2)^{1/4} \quad (8)$$

to get $\Omega_{\tilde{\gamma}} = 1$. However, in most models $m_{\tilde{\chi}}$ and $m_{\tilde{q}}$ are renormalized differently by interactions at scales $\gtrsim m_W$, and a universal value m_0 at some high GUT scale yields different masses. For example, in a minimal supergravity model (MSM), one finds²⁰⁾

$$m_{\tilde{q}}^2 \simeq m_0^2 (1 + 6.4 \xi^2), m_{\tilde{e}_R}^2 \simeq m_0^2 (1 + 0.15 \xi^2), m_{\tilde{e}_L}^2 \simeq m_0^2 (1 + 0.52 \xi^2) \quad (9)$$

where $\xi \equiv m_{1/2}/m_0$ and

$$m_{\tilde{g}} \simeq 0.47 \xi m_0, m_{\tilde{g}} \simeq 2.8 \xi m_0 \quad (10)$$

Assuming these mass relations, we use Eqs. (8) and (11) of Ref. 12) to fix the values of $m_{\tilde{\gamma}}$ needed to obtain any specific value of $\Omega_{\tilde{\gamma}}$. One needs masses given approximately²⁰⁾ by

$$\text{MSM}_1 : m_{\tilde{q}} \simeq 44 \text{ GeV} + 5.5 m_{\tilde{g}} \quad (11)$$

to get $\Omega_{\tilde{\gamma}} = 1$, and²⁰⁾

$$\text{MSM}_{1/4} : m_{\tilde{q}} \simeq 24 \text{ GeV} + 5.3 m_{\tilde{g}} \quad (12)$$

to get $\Omega_{\tilde{\gamma}} = 1/4$. We will use EM(8), MSM_1 (11) and $\text{MSM}_{1/4}$ (12) as examples for the subsequent extraction of model-dependent upper bounds on $\rho_{\tilde{\gamma}}$ and bounds on $m_{\tilde{\gamma}}$.

We now consider uncertainties in the matrix element factors Δq (7). Neutron decay tells us that

$$g_A = \Delta u - \Delta d = 1.25 \quad (13)$$

The naïve quark model is often used^{4),6),11)-14),21)} to set $\Delta s \simeq 0$ and fix Δu and Δd separately. An improved estimate can be made by using flavour SU(3) to determine Δu and Δd by relating them to the F and D parameters measured in hyperon β -decays²²⁾:

^{*}) We assume a present-day Hubble expansion rate $H_0 = 50 \text{ km/s/Mpc}$. If H_0 is larger (smaller), one needs larger (smaller) values of $m_{\tilde{q}}$.

$$\Delta u - \Delta d = F + D, \quad \Delta u + \Delta d - 2\Delta s = 3F - D \quad (14)$$

where $F/D = 0.63^{**}$, and it is again conventionally assumed²³⁾ that $\Delta s \approx 0$ yielding

$$H_B : \quad \Delta u = 0.96(6), \quad \Delta d = -0.28(4) \quad [\Delta s = 0] \quad (15)$$

Recent EMC data²⁵⁾ on polarized μp scattering cast doubt on the assumption that $\Delta s \approx 0$. They measure the spin-dependent structure function $g_1^p(x, Q^2)$, for which there is a QCD sum rule:

$$\int_0^1 dx g_1^p(x, Q^2) = \frac{1}{2} \left(\sum_q Q_q^2 \Delta q \right) (1 + O(\frac{\alpha_s}{\pi})) \approx \frac{1}{2} \left(\frac{4}{9} \Delta u + \frac{1}{9} \Delta d + \frac{1}{9} \Delta s \right) \quad (16)$$

The $O(\alpha_s/\pi)$ correction in (16) is $(-1 \cdot \alpha_s/\pi)$ for non-singlet combinations³³⁾ of the Δq , and $(-1/3 \cdot \alpha_s/\pi)$ for the singlet combination³⁴⁾: these do not give large Q^2 -dependence for $Q^2 > 3 \text{ GeV}^2$. Note that EMC measures the same combination of the Δq as appears in $\sigma(\tilde{\gamma}p \rightarrow \tilde{\gamma}p)$ (7a) if all the m_q are approximately equal as we expect. The prediction using H_B (15) is

$$\int_0^1 dx g_1^p(x, Q^2) = 0.20 \approx \frac{1}{2} \left(\frac{4}{9} \Delta u + \frac{1}{9} \Delta d + \frac{1}{9} \Delta s \right) \quad (17)$$

whereas the EMC measure²⁵⁾

$$\int_0^1 dx g_1^p(x, Q^2) \Big|_{Q^2 \approx 3 \text{ GeV}^2} = 0.113 \pm 0.012 \pm 0.025 \quad (18)$$

If one neglects any possible dependence on the renormalization scale $\mu \approx Q$, this can be combined with (14) to yield

$$EMC : \quad \Delta u = 0.70(8), \quad \Delta d = -0.54(2), \quad \Delta s = -0.25(7) \quad (19)$$

each with errors ± 0.08 (4) if the statistical and systematic errors in (18) are added in quadrature^{**)}.

Different estimates of the factor $(\sum Q_q^2 \Delta q)^2$ appearing in Eq. (7a) are shown in Table 1. The reduction in the observed value (18) compared to that previously

*) A previous estimate²¹⁾ used an out-of-date value for F/D .

**) The estimates (19) correspond to a fraction $-0.09(1) \pm 0.24(2)$ of the proton spin being carried by quarks. They also correspond to a singlet/octet matrix element ratio of $-0.13(3) \pm 0.37(1)$ times the quark model value previously estimated assuming $\Delta s = 0$.

expected (17) suggests a reduction in the prediction of $\sigma(\tilde{\gamma}p \rightarrow \tilde{\gamma}p)$ by a factor ~ 4 . As already mentioned, we do not expect significant renormalization scale (μ) dependence that decouples the EMC measurements at $\mu^2 \approx Q^2 > 3 \text{ GeV}^2$ from the values of the Δq at $\mu^2 \approx m_q^2 \sim 10^4 \text{ GeV}^2$ required in $\tilde{\gamma}p$ scattering. Moreover, the EMC Collaboration sees no evidence for any Q^2 -dependence^{*)}. Until the EMC result is confirmed, one should nevertheless perhaps regard the range (15), (19) as a legitimate uncertainty in dark matter detection estimates.

We now consider the possibility that the dark matter particle is a Higgsino $\tilde{h}$, for which²¹⁾

$$\sigma(\tilde{h}p \rightarrow \tilde{h}p) = \frac{6m_{\tilde{h}}m_p^2}{\pi(m_{\tilde{h}}+m_p)^2} G_F^2 \cos^2(2\beta) \left[\frac{2}{9} I_{3q} \Delta q \right]^2 \quad (20)$$

where $I_{3u} = +\frac{1}{2}$, $I_{3d,s} = -\frac{1}{2}$, Δq is given by (7b), (15) and (19) as before, and β is given by $\tan\beta \equiv v_2/v_1$ where $v_2(v_1)$ is the v.e.v. of the Higgs giving masses to the $u(d,s)$ quarks. Notice that the squark masses do not affect $\sigma(\tilde{h}\tilde{h} \rightarrow f\bar{f})$ or $\sigma(\tilde{h}p \rightarrow \tilde{h}p)$ (20), but it is possible²⁾ to find an acceptable choice of β which gives $\Omega_{\tilde{h}} = 1$ whenever $m_{\tilde{h}} > m_1 \approx 6 \text{ GeV}$.

The factors a_i in (1) are simple numerical constants in νN scattering cross-sections. For n and p targets they are¹⁵⁾

$$a_{\nu p} = 4.2, \quad a_{\bar{\nu}p} = 4.1; \quad a_{\nu n} = 9.5, \quad a_{\bar{\nu}n} = 2.2 \quad (21)$$

corresponding to

$$\left. \begin{array}{ll} \text{isoscalar target: } a_{\nu} = 6.85, & a_{\bar{\nu}} = 3.15 \\ \text{Fe} = 26p + 30n : a_{\nu} = 7.04, & a_{\bar{\nu}} = 3.09 \\ \text{H}_2\text{O} = 10p + 8n : a_{\nu} = 6.56, & a_{\bar{\nu}} = 3.26 \end{array} \right\} \quad (22)$$

The factors B_f in (1) are the branching ratios for $\tilde{X}\tilde{X} \rightarrow f\bar{f}$ which are determined by

^{*)} Jaffe²⁶⁾ has pointed out that there could in principle be significant μ -dependence between $\mu^2 = 3 \text{ GeV}^2$ and $\mu^2 \ll 1 \text{ GeV}^2$. This is irrelevant to $\sigma(\tilde{\gamma}p \rightarrow \tilde{\gamma}p)$, but could invalidate the derivation of the individual Δq in (19). This introduces an uncertainty in the subsequent estimate of $\sigma(\tilde{h}p \rightarrow \tilde{h}p)$, which is fortunately not large.

the couplings in the particular model studied. We use the Monte Carlo procedure detailed in Ref. 15) to model the annihilation events in the Sun and calculate the factors $\langle Nz \rangle_{fi}$ in (1). Effects due to the solar medium are important. Relativistic light hadrons and muons are stopped before they decay, so that only the semi-leptonic decays of heavy quarks and τ decays contribute to the high-energy neutrino flux^{11),15)}. In general, one might also worry about neutrino interactions with the Sun and dE/dx effects on heavy hadrons, but these are negligible for the range of particle masses considered here¹⁵⁾. Each annihilation event gives two fundamental fermions in the final state with energies equal to $m_{\tilde{X}}$. If the annihilation is into a pair of quarks, fragmentation and other final-state QCD effects are important. Heavy-flavour hadrons will not have the "beam" energy $m_{\tilde{X}}$, and the subsequent neutrino spectra are softer than one would naïvely expect. This fragmentation effect has been modelled carefully¹⁵⁾ using a Monte Carlo³⁵⁾ which reproduces correctly the spectra of charged leptons and other secondary particles in e^+e^- annihilation. It reduces $\langle Nz \rangle_{fi}$ by a factor ~ 3 for the case $f = c$, and ~ 2 for the case $f = b$. We will present and use results for $i = \nu_e, \nu_\mu, \bar{\nu}_e, \bar{\nu}_\mu$ only.

Our results are shown in Table 2 and the figures. Table 2 makes predictions for contained high-energy solar neutrino event rates R_ν for the following different models:

- (A) $\tilde{X} = \tilde{\gamma} : \Omega_{\tilde{\gamma}} = 1$, MSM₁ masses and H β estimates of Δq ,
- (B) $\tilde{X} = \tilde{\gamma} : \Omega_{\tilde{\gamma}} = 1$, MSM₁ masses and EMC estimates of Δq ,
- (C) $\tilde{X} = \tilde{\gamma} : \Omega_{\tilde{\gamma}} = 1$, EM₁ masses $m_{\tilde{X}} = m_{\tilde{q}}$ and H β estimates of Δq ,
- (D) $\tilde{X} = \tilde{\gamma} : \Omega_{\tilde{\gamma}} = 1/4$, MSM_{1/4} masses and H β estimates of Δq ,
- (E) $\tilde{X} = \tilde{h} : \Omega_{\tilde{h}} = 1$, H β estimates of Δq ,
- (F) $\tilde{X} = \tilde{h} : \Omega_{\tilde{h}} = 1$, EMC estimates of Δq .

We see that only in case (D) does the rate predicted for $\rho_{\tilde{X}} = 0.2 \text{ GeV cm}^{-3}$ and $\bar{v} = 300 \text{ kms}^{-1}$ come above the upper limit of $R_\nu < 1.6 \text{ events kton}^{-1}\text{y}^{-1}$ provided by the Fréjus experiment¹⁸⁾. An upper limit on contained events is also available from the IMB experiment^{16),*}. The Kamioka experiment is also preparing an upper limit on R_ν ¹⁷⁾, but this is not yet available in a form which can be compared with the astrophysically expected¹⁹⁾ halo density (6a). Fig. 1a shows upper limits from the Fréjus experiment¹⁸⁾ on $\rho_{\tilde{\gamma}}$ in cases (A), (B) and (C) above. Also shown is the

*) The IMB experiment¹⁶⁾ also quotes an upper limit on the upward-going μ flux, but this is not yet competitive with upper limits on the ν flux from contained events at Fréjus.

upper limit from the IMB experiment¹⁶⁾ on $\rho_{\tilde{\gamma}}$ in case (A). The horizontal lines indicate the expected range of the halo density¹⁹⁾, which is not yet excluded by the underground experiments for any value of $m_{\tilde{\gamma}}^*$. Figure 1b shows upper limits from the Fréjus experiment¹⁸⁾ on $\rho_{\tilde{h}}$ in cases (E) and (F) above, where again we see that no region of $m_{\tilde{h}}$ can be excluded in a model-independent way. Since the expected sparticle masses decrease when $\Omega_{\tilde{\gamma}}$ is decreased [see (11) and (12)], resulting in an increase in $\sigma(\tilde{\gamma}p \rightarrow \tilde{\gamma}p)$ (7a) and hence an increase in $R_{\tilde{\gamma}}$ (1), we might expect to get a non-trivial bound on $m_{\tilde{\gamma}}$ if $\Omega_{\tilde{\gamma}}$ is sufficiently small. This effect was already discernible in case (D) of Table 2, where $\Omega_{\tilde{\gamma}} = 1/4$ and the H β estimates of Δq are used. It is also exhibited in Fig. 2, as a domain of the $(m_{\tilde{\gamma}}, \Omega_{\tilde{\gamma}})$ plane which can be excluded assuming MSM masses (9), (10) and H β estimates of Δq (15). Note, however, that even this exclusion zone is greatly reduced if the EMC estimates of Δq (19) are used instead.

So far, we have concentrated on the contained event signal for high-energy solar neutrinos. Alternatively, one can look for upward-going muons from ν_{μ} collisions in rock^{6), 14)}. With the same assumptions and uncertainties as in the derivation of Eq. (1), we estimate

$$\begin{aligned}
 R_{\mu} \approx & 1.1 \times 10^{-2} \left(\frac{\text{Area}}{100 \text{ m}^2} \right) \left(\frac{\rho_L}{0.3 \text{ GeV cm}^{-3}} \right) \left(\frac{300 \text{ km s}^{-1}}{v} \right) \times \\
 & \times \Theta(m_{\tilde{\gamma}} - m_E) f(m_{\tilde{\gamma}}/m_c) \left(\frac{\sigma(\tilde{\gamma}p \rightarrow \tilde{\gamma}p)}{10^{-40} \text{ cm}^2} \right) \\
 & \times \sum_i a_i b_i \sum_f B_f \langle N z^2 \rangle_{fi} \quad : \quad v_{\mu} = 0.5, v_{\tilde{\gamma}} = 0.67
 \end{aligned} \tag{23}$$

for the rate of such through-going muon events. The existing Fréjus¹⁸⁾, IMB¹⁶⁾ and Kamioka¹⁷⁾ experiments are unlikely to improve their current limits on $R_{\tilde{\gamma}}$ (1) by large factors, and the IMB limit¹⁶⁾ on R_{μ} (23) is not yet competitive. However, future experiments with large underground μ detectors such as MACRO³⁶⁾ could make significant improvements in the upper limit on R_{μ} . Therefore, we have included in Table 2 predictions for R_{μ} in the models (A) and (E) introduced above. This is likely to be in the future the best way to establish lower bounds on $m_{\tilde{\gamma}}$ and $m_{\tilde{h}}$ or to find indirect evidence for supersymmetric dark matter particles.

*) The vertical glitches in these curves correspond to $\tilde{\gamma}\tilde{\gamma}$ annihilations in the early Universe at temperatures $T_A \sim 1/20 m_{\tilde{\gamma}}$ higher/lower than the hadron phase transition temperature, taken here to be $T_c = 300 \text{ MeV}$. The values of $m_{\tilde{\gamma}}$ required to get $\Omega_{\tilde{\gamma}} = 1$ are different for $T_A \gtrless T_c$, causing a change in $\sigma(\tilde{\gamma}p \rightarrow \tilde{\gamma}p)$. Because T_c is uncertain, so also are the exact values of $m_{\tilde{\gamma}}$ where the glitches occur.

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Table 1

Estimates of $\bar{M}^2 \equiv (\sum_q Q_q^2 \Delta q)^2$

References	4),11)	21)	This paper, using 23)	This paper, using 25)
Assumptions	$\tilde{u}$ only; naïve quark model	$\tilde{u} + \tilde{d};$ $\Delta s = 0,$ old F/D	$\tilde{u} + \tilde{d};$ $\Delta s = 0,$ F/D = 0.63	$\tilde{u} + \tilde{d} + \tilde{s};$ EMC,
for protons	0.24	0.11	0.16	0.05
for neutrons	0.005	0.007	0.0004	0.04

Table 2

Model-dependent estimates of R_ν , R_μ . These numbers and the Fréjus curves in the figures are calculated for $E_\nu > 2 \text{ GeV}^{18)}$. The IMB curves in the figures are for $0.2 < E_\nu < 2 \text{ GeV}^{16)}$.

	R_ν (events/kty)						R_μ (events/100m ² y)	
Model →	A	B	C	D	E	F	A	E
	$\tilde{\gamma}$				$\tilde{h}$		$\tilde{\gamma}$	$\tilde{h}$
$m_{\tilde{X}}$ (GeV)	$\Omega=1$	$\Omega=1$	$\Omega=1$	$\Omega=1/4$	$\Omega=1$	$\Omega=1$	$\Omega=1$	$\Omega=1$
	MSM ₁	MSM ₁	EM ₁	MSM _{1/4}	H β	EMC	MSM ₁	H β
↓	H β	EMC	H β	H β			H β	
3	0.46	0.15	0.51	1.91	—	—	0.12	—
4	1.56	0.50	1.80	6.10	—	—	0.44	—
6	1.83	0.59	2.28	6.24	1.16	1.69	0.68	0.35
8	0.82	0.26	1.03	2.74	0.61	0.89	0.37	0.21
10	0.58	0.19	0.81	1.76	0.59	0.86	0.31	0.23
15	0.22	0.072	0.39	0.57	0.33	0.49	0.17	0.17
20	0.11	0.034	0.23	0.23	0.22	0.32	0.11	0.14
30	0.028	0.009	0.09	0.05	0.10	0.15	0.04	0.09

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FIGURE CAPTIONS

Fig. 1 : Upper limits on $\rho_{\tilde{X}}$: (a) $\tilde{X} = \tilde{\gamma}$, using Fréjus data¹⁸⁾ under model assumptions (A), (B) and (C), and IMB data¹⁶⁾ under assumption (A); (b) $\tilde{X} = \tilde{h}$ using Fréjus data¹⁸⁾ under model assumptions (E), (F). The $\hat{E}$ and $\hat{F}$ curves are obtained from Eq. (20) by removing the model-dependent factor of $\cos^2(2\beta)$. The horizontal lines indicate the astrophysically preferred value and probable limits¹⁹⁾ on $\rho_{\tilde{X}}$.

Fig. 2 : Domain of the $(m_{\tilde{\gamma}}, \Omega_{\tilde{\gamma}})$ plane excluded by the IMB¹⁶⁾ and Fréjus data¹⁸⁾ assuming MSM masses (9) and (10) and the H β and EMC estimates of Δq (15).

Fig. 1a

Fig. 1b

Fig. 2