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SUPERSTRING DARK MATTER

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ABSTRACT

We argue that the lightest supersymmetric particle (LSP) emerging from the superstring theory is a mixture of neutral gauginos and matter fermions. We present their mixing matrix calculated in a plausible minimal low-energy model abstracted from the superstring and exhibit the composition of the LSP χ . We compute its relic cosmological density and find that it lies within a factor 2 of the critical density required for closure, over a wide range of possible input parameters. We also compute the flux of neutrinos from LSP annihilation in the Sun, and find that it straddles the upper bound from proton decay detectors. Acceptable fluxes are obtained if $m_\chi < m_t$, in which case the superstring relic can have the critical density for a present Hubble expansion rate $H_0 \gtrsim 50$ km/s/Mpc only if $m_t \gtrsim 40$ GeV.

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We cannot see most of the matter in the Universe, and would like to understand the nature of this dark matter. One of the most attractive hypotheses is that it is some species of stable, massive, weakly-interacting elementary particle. Just such a particle is predicted by most supersymmetric models, where the lightest supersymmetric particle (LSP) is usually stable because of R parity conservation¹⁾. Many models for the dark matter invoke mass scales which are very different from that of conventional baryonic matter²⁾, and hence do not naturally predict a dark matter density ρ_D similar to that of baryons (ρ_B). Since the baryon density ρ_B is within two orders of magnitude³⁾ of the critical density ρ_C required for closure, this means that such models do not naturally predict a dark matter density ρ_D similar to the closure density ρ_C , and one must tune parameters to force $\rho_D \sim 0(\rho_C)$ as required by observation and by inflation. This objection need not apply to supersymmetric models^{4),5)}, whose mass scale of order 100 GeV is not much different from that of baryons. Indeed, calculations⁶⁾ in the minimal supersymmetric extension of the Standard Model gave a lower bound $\rho_D \gtrsim 10^{-2}\rho_C$ and could naturally accommodate $\rho_D \approx \rho_C$.

Nowadays the minimal supersymmetric Standard Model has been displaced in the affections of theorists by the $E_8 \times E_8'$ superstring⁷⁾, and one inevitably asks what dark matter candidates it contains, and what their relic density ρ_D might be. The $E_8 \times E_8'$ superstring contains two gauge sectors which are linked only by interactions of gravitational strength. The hidden sector contains no matter fields, only gauge bosons and gauginos⁸⁾. The hidden sector gauge interactions are expected to become strong at some high mass scale and have a spectrum of bound states which are gauge and flavour singlets. There is therefore no quantum number to prevent them from decaying into pairs of observable sector particles, and one does not expect the hidden sector to contain a stable dark matter candidate. On the other hand, the observable sector does contain particles of spin 0 and of spin $\frac{1}{2}$ which could be the LSP and provide the dark matter, much as the sneutrino⁹⁾ and gaugino/shiggs mixture⁵⁾ could in the minimal supersymmetric Standard Model. Just as in that model, the generic tendency is that one of the spin- $\frac{1}{2}$ particles weighs less than any of the spin-0 sparticles, so that the most likely superstring candidate for dark matter is some combination of the weakly-interacting gauginos and matter fermions in the observable sector.

In this paper, we analyze the cosmological relic density expected for such a superstring LSP, and show that $\Omega_D \equiv \rho_D/\rho_C \sim 1$ over a wide range of the plausible

parameter space in a minimal model^{10),11)} inspired by the superstring^{*}). The model automatically gives us a density of dark matter which is useful for the formation of structure in the Universe, as well as being required by inflation. We also calculate the rates of capture¹²⁾ and annihilation¹³⁾ of these superstring LSPs in the Sun, and show that they give a flux of high-energy neutrinos which is comparable with the experimental limits from proton decay detectors¹⁴⁾. Thus, superstring dark matter may be detectable as well as inevitable.

First we analyze the composition of the superstring LSP which is our dark matter candidate. After Calabi-Yau compactification, matter fermions from the superstring¹⁰⁾ occupy 27-dimensional representations of some subgroup of E_6 which have the following $SO(10)$ and $SU(5)$ decompositions:

$$\begin{aligned} \underline{27} &= \underline{16} + \underline{10} + \underline{1}(\tilde{N}) \\ &= (\underline{10} + \underline{5} + \underline{1}(\nu^c)) + (\underline{5} + \underline{5}) + \underline{1}(\tilde{N}) \end{aligned} \quad (1)$$

Indicated in parentheses are the two novel weakly-interacting neutrals ν^c and $\tilde{N}$ which appear in each 27-dimensional generation along with neutral particles from the $SO(10)$ $\underline{10} = \underline{5} + \underline{5}$ of $SU(5)$, which have the quantum numbers of conventional Higgs doublets H and $\bar{H}$. Thus we have at least three generations of $(\tilde{H}^0, \tilde{\bar{H}}^0, \tilde{N}, \nu, \nu^c)$ neutral matter fermions. If R parity is conserved as we expect, the $(\tilde{H}^0, \tilde{\bar{H}}^0, \tilde{N})$ which have $R = -1$ can mix with the gauginos to form the LSP, but not with the ν and ν^c which have $R = +1$. There is a non-trivial constraint on the parameters of the model which must be obeyed if the additional light ν^c particles are not to affect unduly primordial nucleosynthesis¹⁵⁾. This constraint is used in our later analysis. The minimal possible gauge group¹⁰⁾ is $SU(3)_C \times SU(2)_L \times U(1)_Y \times U(1)_E$, where the $U(1)_E$ quantum numbers of $\tilde{H}^0$, $\tilde{\bar{H}}^0$ and $\tilde{N}$ are $-4/3$, $-1/3$ and $5/3$ respectively, with equal coupling constants for the two hypercharges. It is possible¹¹⁾ to construct viable dynamical models based on this minimal gauge group, and we do not know how to construct models based on non-minimal gauge groups. Therefore we will not consider them in detail in what follows, though they should share the qualitative features of our analysis. Thus the neutral supersymmetric fermions which we believe mix to form the LSP are $(\tilde{W}_3, \tilde{B}, \tilde{H}^0, \tilde{\bar{H}}^0, \tilde{B}_E, \tilde{N})$. In principle, we should consider the possibility that each one of $N_g = 3$

*) We also give arguments that this result is likely also in a range of other superstring-inspired models.

generations of the matter fermions ($\tilde{H}^0$, $\tilde{H}^\pm$, $\tilde{N}$) could participate in the mixing, but we will first make the simplifying assumption that only one generation need be considered. This restriction will be justified a posteriori. The full (6×6) mixing matrix we consider is:

$$\begin{pmatrix}
 (\tilde{W}^3, \tilde{B}, \tilde{H}^0, \tilde{H}^\pm, \tilde{B}_E, \tilde{N}) / M_2 & 0 & -\frac{g_2 v}{\sqrt{2}} & \frac{g_2 \bar{v}}{\sqrt{2}} & 0 & 0 \\
 & \frac{5}{3} \frac{\alpha_1}{\alpha_2} M_2 & \frac{g_1 v}{\sqrt{2}} & -\frac{g_1 \bar{v}}{\sqrt{2}} & 0 & 0 \\
 & & 0 & \lambda x & \frac{2\sqrt{2}}{3} g_1 v & \lambda \bar{v} \\
 & & & 0 & -\frac{1}{3\sqrt{2}} g_1 \bar{v} & \lambda v \\
 & & & & \frac{5}{3} \frac{\alpha_1}{\alpha_2} M_2 & \frac{5}{3\sqrt{2}} g_1 x \\
 & & & & & 0
 \end{pmatrix}
 \begin{pmatrix}
 \tilde{W}^3 \\
 \tilde{B} \\
 \tilde{H}^0 \\
 \tilde{H}^\pm \\
 \tilde{B}_E \\
 \tilde{N}
 \end{pmatrix} \quad (2)$$

symmetric

where M_2 is the effective SU(2) gaugino mass, we have incorporated the corresponding renormalized values for the U(1) gaugino masses, $v \equiv \langle 0 | H | 0 \rangle$, $\bar{v} \equiv \langle 0 | \bar{H} | 0 \rangle$, $x \equiv \langle 0 | N | 0 \rangle$ and λ is the $\overline{H}HN$ superpotential coupling. The mass eigenstates of the matrix (2) are six Majorana fermions χ_i with the general decomposition

$$\chi_i = \alpha_i \tilde{W}^3 + \beta_i \tilde{B} + \gamma_i \tilde{H}^0 + \delta_i \tilde{H}^\pm + \epsilon_i \tilde{B}_E + \phi_i \tilde{N} \quad (3)$$

To study the matrix (2) numerically, we need input values for the parameters (M_2 , v , $\bar{v}$, x , λ), which we take from some dynamical calculations¹¹⁾ in this minimal superstring model. We assume¹¹⁾ that the dominant source of supersymmetry breaking in the observable sector at the unification scale is the gaugino mass $m_{\frac{1}{2}}$, which yields the following soft supersymmetry-breaking masses after renormalization corrections:

$$m_{\frac{1}{2}} : M_2 : m_{\tilde{t}_L} : m_{\tilde{t}_R} : m_{\tilde{g}} = 1 : 0.33 : 0.70 : 0.38 : 1.95 \quad (4)$$

We do not assume any fixed value of $m_{\frac{1}{2}}$, but use the following empirical correlation¹¹⁾ between x/v and $m_{\frac{1}{2}}$:

$$\frac{x}{v} = \frac{(m_{1/2} + 75 \text{ GeV})}{88 \text{ GeV}} : M_2 = 0.33 m_{1/2} \quad (5)$$

We have considered the range of values

$$1.7 \leq \frac{v}{v} \leq 5 \quad (6)$$

which was favoured in the dynamical calculations, and different values of λ in the plausible range^{*)}

$$0.10 \leq \lambda \leq 0.25 \quad (7)$$

The results for $v/\bar{v} = 2$, $\lambda = 0.15$ are shown in Fig. 1: other values of the parameters in the ranges (6) and (7) give similar results. We see that for $m_{1/2} > 100$ GeV, as suggested by phenomenological particle physics constraints¹¹⁾, notably the requirement that $m_{\tilde{v}} > 0$ after the inclusion of D terms, the LSP is mainly composed of the Standard Model fields ($\tilde{W}^3$, $\tilde{B}$, $\tilde{S} \equiv (\tilde{v}\tilde{H}^0 + v\tilde{H}^0)/\sqrt{v^2 + \bar{v}^2}$, $\tilde{A} \equiv (v\tilde{H}^0 - \bar{v}\tilde{H}^0)/\sqrt{v^2 + \bar{v}^2}$) with relatively small admixtures of the novel fields ($\tilde{B}_E$, $\tilde{N}$).

We now turn to the calculation of the relic cosmological mass density of the LSP. A previous general analysis⁵⁾ tells us that it is

$$\rho_{\tilde{\chi}} = 4.9 \times 10^{-40} \text{ g/cm}^3 \times \left(\frac{T_{\tilde{\chi}}}{T_{\gamma}} \right)^3 \left(\frac{T_{\gamma}}{2.8^{\circ}\text{K}} \right)^3 N_F^{1/2} \left(\frac{1 \text{ GeV}^{-2}}{\tilde{a} x_f + \frac{1}{2} \tilde{b} x_f^2} \right) \quad (8)$$

where $T_{\gamma} \approx 2.8^{\circ}\text{K}$ is the apparent temperature of the microwave background radiation today, and $T_{\tilde{\chi}}$ is the equivalent relic temperature which depends on its freeze-out temperature $T_f \equiv x_f m_{\tilde{\chi}}$ as displayed in Ref. 16). The quantities $\tilde{a}$ and $\tilde{b}$ parametrize the thermally-averaged annihilation cross-section:

$$\langle \sigma v_{\text{rel}} \rangle \equiv \tilde{a} + \tilde{b} x \quad (9)$$

where $\langle v_{\text{rel}}^2 \rangle = 6x$: $x \equiv T/m_{\tilde{\chi}}$. They are related to the effective interaction parameters A_f , B_f [here $L, R \equiv (1 \mp \gamma_5)/2$]:

*) If one took $\lambda < 0$, minimization of the effective scalar potential including the soft supersymmetry-breaking trilinear term proportional to $\lambda x v \bar{v}$ would force us to take $\bar{v} < 0$ also, resulting in a mass matrix equivalent to (2).

$$\mathcal{L}_{f\bar{f}} = (\bar{\chi} \gamma^\mu \gamma_5 \chi) (\bar{f} \gamma_\mu (A_f L + B_f R) f) \quad (10)$$

by

$$\begin{aligned} \tilde{a} &= \sum_f \frac{\Theta(m_\chi - m_f)}{2\pi} \sqrt{1 - \frac{m_f^2}{m_\chi^2}} m_f^2 (A_f - B_f)^2 \\ \tilde{b} &= \sum_f \frac{\Theta(m_\chi - m_f)}{2\pi} \sqrt{1 - \frac{m_f^2}{m_\chi^2}} \left\{ (A_f + B_f)^2 \left(4m_\chi^2 - \frac{13}{4} m_f^2 \right) \right. \\ &\quad \left. + \frac{21}{2} A_f B_f m_f^2 + \frac{3}{4} \frac{m_f^2 (A_f - B_f)^2}{(1 - m_f^2/m_\chi^2)} \right\} \end{aligned} \quad (11)$$

In our model with the LSP χ parametrized by the mixing parameters $(\alpha, \beta, \gamma, \delta, \epsilon, \phi)$ of Eq. (3), the parameters A_f, B_f are given by

$$A_f = A_f^{\tilde{f}} + A_f^Z + A_f^E, \quad B_f = B_f^{\tilde{f}} + B_f^Z + B_f^E \quad (12)$$

where we have separated out the $\tilde{f}$, Z and Z_E exchange contributions. They are given by

$$\begin{aligned} A_f^{\tilde{f}} &= \frac{(T_{f_L}^3 \alpha g_2 + \frac{1}{2} Y_{f_L} \beta g_1 + \frac{1}{2} Y_{f_L}^E \epsilon g_E)^2}{2m_{\tilde{f}_L}^2} - \begin{cases} \gamma^2 m_f^2 / 4v^2 m_{\tilde{f}_R}^2 & (f = u, c, t) \\ \delta^2 m_f^2 / 4v^2 m_{\tilde{f}_R}^2 & (f = d, s, b, \\ & e, \mu, \tau) \end{cases} \\ B_f^{\tilde{f}} &= -A_f^{\tilde{f}} (f_L \leftrightarrow f_R) \end{aligned} \quad (13)$$

where the T^3 , Y and Y^E quantum number assignments for light fermions are shown in the Table, $g_E = g_1 = \tan\theta_w g_2$,

$$A_f^Z = (\gamma^2 - \delta^2) \left(\frac{g_2 \cos \theta_W + g_1 \sin \theta_W}{4m_Z^2} \right) \left(\frac{Y_{f_L}}{2} g_1 \sin \theta_W - T_{f_L}^3 g_2 \cos \theta_W \right)$$

$$B_f^Z = A_f^Z (f_L \leftrightarrow f_R) \quad (14)$$

and

$$A_f^E = \left(\frac{5}{6} \phi^2 - \frac{2}{3} \gamma^2 - \frac{1}{6} \delta^2 \right) \frac{g_E^2}{2m_{Z_E}^2} \left(\frac{Y_{f_L}^E}{2} \right)$$

$$B_f^E = A_f^E (f_L \leftrightarrow f_R) \quad (15)$$

In writing the expressions (12), (14) and (15), we have neglected mixing between the conventional $SU(2)_L \times U(1)_Y$ neutral boson Z and the additional $U(1)_E$ gauge boson Z_E . In fact, phenomenological constraints, notably¹¹⁾ the absence of a significant mass-shift for the observed Z^0 boson [and more stringently¹⁵⁾ if less certainly the absence of an effect on primordial nucleosynthesis due to the ν^c interacting via Z_E exchange if the ν^c is light], tell us that $x/v \gg 1$, $m_{Z_E} \gg m_Z$, and mixing is small. In what follows, we can therefore neglect the Z_E contributions (A_f^E , B_f^E) to superstring LSP annihilation. Thus the only change from our⁵⁾ previous formulae in a minimal supersymmetric Standard Model is the ϵ term in Eq. (13), which has a small effect because ϵ is generally small, as seen in Fig. 1. Similar approximations would hold in other low-energy models motivated by the superstring, so our subsequent results should have more general validity.

Numerical results using Eqs. (8) to (14) are shown in Figs. 2 and 3. They have been obtained assuming $m_t = 40$ GeV and using the relation (4) between the soft supersymmetry-breaking parameters M_2 , $m_{\tilde{q}}$, $m_{\tilde{\chi}}$ and $m_{\frac{1}{2}}$. We have also included the effects of the D terms¹¹⁾ on the squark and slepton masses, which are somewhat important even for $m_{\frac{1}{2}} > 100$ GeV as considered here^{*)}. We see from Figs. 2 and 3 that in much of the $(\lambda, m_{\frac{1}{2}})$ plane

*) There is a very thin sliver of $m_{\frac{1}{2}}$ values close to this limit where the sneutrino is lighter than the lightest of two supersymmetric fermions. This exceptional case is astrophysically excluded¹⁷⁾, except for unnaturally small values of $m_{\tilde{\nu}}$.

$$\Omega_D \equiv \rho_D / \rho_c = 2^{0 \pm 1} (0.5/h_0)^2 \quad (16)$$

where $\rho_c = 1.88 \times 10^{-29} h_0^2$ gm/cc is the closure density, and the present Hubble expansion rate $H_0 \equiv 100 h_0$ km/s/Mpc. This range of Ω_D is narrower than that allowed in our previous analysis^{5),6)} of a minimal supersymmetric Standard Model, largely because of the restrictive relationships (4) between sparticle masses, which supplant the previous two-parameter set of mass formulae depending on m_0 and $m_{1/2}$. Although it was natural for the superstring dark matter density to be close to ρ_c , as argued in the introduction of this paper, it could have come out at least an order of magnitude wrong. We find the result (16) remarkable, and it encourages us to think that we are on the right track!

Now it is convenient to discuss the incorporation of more generations of $\tilde{H}^\circ$, $\tilde{\bar{H}}^\circ$ and $\tilde{N}$ fermions. It is always possible¹⁸⁾ to work in a field basis where only one each of the H° , $\bar{H}^\circ$ and N have vacuum expectation values v , $\bar{v}$ and x respectively, while the others (denoted by h° , $\bar{h}^\circ$ and n) have zero vacuum expectation values¹⁸⁾. It is easy to check that in this basis there is no mixing between the gauginos and the spartners $\tilde{h}^\circ$, $\tilde{\bar{h}}^\circ$ and $\tilde{n}$ of the scalars without vacuum expectation values. Nor do the $(\tilde{h}^\circ, \tilde{\bar{h}}^\circ, \tilde{n})$ mix with the $(\tilde{H}^\circ, \tilde{\bar{H}}^\circ, \tilde{N})$ in this basis, because such a mixing could only arise via $H\tilde{H}n$, $H\tilde{H}N$ and $h\tilde{H}N$ terms in the superpotential whose associated soft supersymmetry-breaking A terms would have generated $\langle 0|n|0\rangle$, $\langle 0|\bar{h}^\circ|0\rangle$, $\langle 0|h^\circ|0\rangle \neq 0$, which we have rotated away by our choice of basis. Therefore the $(\tilde{W}^3, \tilde{B}, \tilde{B}_E, \tilde{H}^\circ, \tilde{\bar{H}}^\circ, \tilde{N})$ and $(\tilde{h}_{1,2}^\circ, \tilde{\bar{h}}_{1,2}^\circ, \tilde{n}_{1,2})$ fermions have independent mass matrices which do not mix. The masses of eigenstates in the $(\tilde{h}^\circ, \tilde{\bar{h}}^\circ, \tilde{n})$ sector depend on x , v , $\bar{v}$ and unknown Yukawa couplings. These could easily be large enough for the lightest $(\tilde{h}^\circ, \tilde{\bar{h}}^\circ, \tilde{n})$ mass eigenstate to be heavier than the lightest mass eigenstate in the submatrix we have discussed. Cosmology requires this to be the case, because if the LSP were some mixture of $\tilde{h}^\circ$, $\tilde{\bar{h}}^\circ$ and $\tilde{n}$ states, and their couplings had similar magnitudes to those of conventional shiggses, they would have a relic cosmological density $\rho_D \gtrsim 0(100) \times \rho_c$. Therefore we retain as our only LSP candidate the lightest mass eigenstate of the matrix (2) previously discussed, which is mainly a mixture of $\tilde{W}^3$ and $\tilde{B}$ with some $\tilde{H}^\circ$ and $\tilde{\bar{H}}^\circ$, and for which the desirable result $\Omega_D \approx 1$ is natural.

As already mentioned, this result is probably also valid in alternative superstring-inspired models^{19),20)} with different additional $U(1)$ gauge groups surviving down to low energies. Such models could differ from ours in the $(\tilde{B}_E, \tilde{N})$

sector of the mixing matrix (2) - which we have already argued is relatively unimportant, in our assumed constraints (5), (6) and (7) on the vacuum expectation values and Yukawa couplings - which we cannot justify in the absence of dynamical calculations in these rival models, in the renormalizations of the soft supersymmetry-breaking parameters (4) - which is unimportant ($\lesssim 5\%$) except for the $\tilde{\chi}_R$, and in the D term contributions to the sparticle masses - which also do not change our results qualitatively. We have studied the specific model of Ref. 20), in which the additional low energy U(1) is that which allows the ν^c to acquire a large Majorana mass - although a likely dynamical mechanism for generating it is not yet available. In this model the $\tilde{\chi}_R$ mass is somewhat larger than in our preferred model, and ρ_D is correspondingly increased though Ω_D remains $O(2)$ in much of the allowed parameter space.

Thus our result (16) is quite stable, and the superstring LSP is an excellent candidate for the dark matter in galactic halos. As they pass through the Sun, LSPs may become trapped¹²⁾ in it if they have a significant elastic scattering cross-section on protons. The rate at which LSPs get trapped is simply estimated by

$$\Gamma_T \approx \pi R_\odot^2 n_\chi \bar{v} \left(\frac{R_c}{R_\odot} \right)^2 \left(\frac{\bar{v}}{v_{esc}} \right)^2 f_\sigma \quad (17)$$

where $n_\chi \equiv \rho_\chi^{local}/m_\chi$, $\bar{v} \approx 300$ km/s is the r.m.s. velocity of a halo LSP near the Sun, the factor $(R_c/R_\odot)^2 (\bar{v}/v_{esc})^2$ accounts for the fact that LSPs with impact parameters larger than $R_\odot = 2G_N M_\odot / v_{esc}^2$ but smaller than $R_c = 2G_N M_\odot / \bar{v}^2$ can also strike the Sun, and f_σ is the probability that an LSP striking the Sun actually gets trapped. The estimate (17) agrees well with the detailed calculations of Press and Spergel¹²⁾ which yield

$$\Gamma_T = \sqrt{\frac{3\pi}{2}} (2G_N R_\odot M_\odot) n_\chi f_\sigma / \bar{v} \quad (18)$$

with $f_\sigma = 0.2 \sigma_{E,36} \times \min(1, 43 \text{ GeV}/m_\chi)$, where $\sigma_{E,36} \equiv \sigma_E / 10^{-36} \text{ cm}^2$ is the elastic cross-section for $\chi + p \rightarrow \chi + p$ scattering. If indeed $\Omega_D \gtrsim 0.2$, one can keep rotational velocities and mass-to-light ratios in agreement with observations by invoking something like a biased galaxy formation scenario²¹⁾. Then, independently of Ω_D as long as it is $\gtrsim 0.05$ ²²⁾, we can infer the local galactic halo

density of LSPs from the rotation curves: $\rho_D^{\text{local}} \approx 0.3 \text{ GeV/cm}^3$ in the solar neighbourhood²³⁾. Numerically, Eq. (18) then yields

$$\Gamma_T = 9 \times 10^{28} \left(\frac{\rho_\chi}{0.3 \text{ GeV/cm}^3} \right) \frac{\sigma_{E,36}}{\bar{v}_{300} m_\chi} S^{-1} \quad (19)$$

where $\bar{v}_{300} = \bar{v}/300 \text{ km s}^{-1}$. The elastic cross-section σ_E can easily be calculated in terms of the quantities A_f and B_f [Eq. (10)]:

$$\sigma_E = (A_d - B_d)^2 \left(\frac{3}{2} - \frac{1}{4} \right)^2 \frac{3}{\pi} \frac{m_\chi^2 m_p^2}{(m_\chi + m_p)^2} \quad (20)$$

where m_p is the proton mass and $\xi \equiv (A_u - B_u)/(A_d - B_d)$. For $m_\chi \gtrsim 6 \text{ GeV}$ ^{24), 17)} as in our model, the solar abundance of LSPs relative to protons is limited by the rate Γ_A of $\chi\chi$ annihilations, so that in equilibrium $\Gamma_A = \frac{1}{2}\Gamma_T$.

These annihilations in the solar interior are a source¹³⁾ of high-energy neutrinos ($E_\nu \gtrsim 300 \text{ MeV}$) which could compete with the atmospheric background¹⁴⁾ produced by cosmic rays, and could be observed in detectors for proton decay. If we define N_{ν_i} ($= N_{\bar{\nu}_i}$) to be the final number of neutrinos (antineutrinos) of type i to be produced per annihilation, the flux of neutrinos and antineutrinos arriving at the Earth is simply

$$\phi_i = \left(\frac{1}{2} \Gamma_T \right) \frac{(N_{\nu_i} + N_{\bar{\nu}_i})}{4\pi R^2} = 32 N_{\nu_i} \left(\frac{\rho_\chi}{0.3 \text{ GeV/cm}^3} \right) \frac{\sigma_{E,36}}{\bar{v}_{300} m_\chi} \text{ cm}^{-2} \text{ s}^{-1} \quad (21)$$

where R is one Astronomical Unit. We take N_{ν_i} from the number of secondary ν_e and ν_μ produced by π and K decay, as well as the primary prompt decay neutrinos²³⁾:

$$N_{\nu_e} \simeq 0.6 \left(\frac{N_\pi}{N_A} \right) + 3.7 \left(\frac{N_K}{N_A} \right) + 7 \left(\frac{N_b}{N_A} \right) + 9.5 \left(\frac{N_t}{N_A} \right)$$

$$N_{\nu_\mu} \simeq 2 N_{\nu_e} \quad (22)$$

with $\langle E_\nu \rangle \approx 0.05 m_\chi$ for most of the neutrinos, while the prompt ones are more energetic. We can estimate the relative branching ratios (N_f/N_A) for annihila-

tions into the different flavours f by using Eqs. (9) and (10) (recall that in the Sun, annihilations occur at low velocities $v \sim 10^{-3}$ so that $x \ll 1$). Thus

$$\frac{N_f}{N_A} \propto n_{c_f} \left(1 - \frac{m_f^2}{m_\chi^2}\right)^{\frac{1}{2}} m_f^2 (A_f - B_f)^2 \quad (23)$$

where n_{c_f} is the number of colours of the flavour f . The fluxes (20) and (21) can now be compared with the atmospheric background which falls off as E_ν^{-2} .

A sampling of our numerical results is shown in Fig. 2. We see that when $\lambda = 0.15$ our solar neutrino fluxes are considerably lower than the atmospheric background if $m_\chi < m_t$, corresponding to $m_{\frac{1}{2}} \lesssim 6 m_t \approx 240$ GeV for our choice of $m_t = 40$ GeV. The small rises in the neutrino fluxes for $m_{\frac{1}{2}} \approx 200$ GeV occur because the LSP is starting to acquire a larger shiggs component which can scatter via Z° exchange, enhancing the trapping rate and also the probabilities (N_f/N_A) for $\chi\chi$ annihilations into heavier flavours. The solar neutrino fluxes in this range of $m_{\frac{1}{2}}$ might eventually be detectable in an experiment with good angular resolution. When $m_\chi \gtrsim m_t$ ($m_{\frac{1}{2}} \gtrsim 240$ GeV in our case), the solar neutrino flux jumps above 10% of the atmospheric background. This could presumably be detected in a dedicated analysis of the available proton decay detector data by exploiting energy and angle information, and we guess that it is unacceptably large, although no experimental upper limit has yet been published. For larger (smaller) values of λ , the shiggs component of the LSP is decreased (increased), resulting in a decrease (increase) in the calculated solar neutrino flux. Values of $\lambda < 0.1$ probably give unacceptably high neutrino fluxes, even if $m_\chi < m_t$. It is striking that the solar neutrino constraint allows ρ_D to creep up to ρ_c only if m_t is sufficiently large: $m_t \ll 40$ GeV would not do. This correlation could be expressed in two remarkable ways. Either one could use a particle physics observation that $m_t \approx 40$ GeV to give an upper bound $m_{\frac{1}{2}} \lesssim 240$ GeV and hence $\Omega_D \lesssim (0.5/h_0)^2$, or one could say that because inflation requires $\Omega_D \approx 1$, cosmology predicts $m_t \gtrsim 40$ GeV if $h_0 = \frac{1}{2}$.

We conclude this paper with a few comments on the possibility of detecting in the laboratory²⁵⁾ the superstring dark matter discussed in this paper. When $m_\chi \sim 20$ GeV, our LSP interacts with matter much like a conventional photino, albeit via the exchange of relatively heavy squarks with masses $m_{\tilde{q}} \sim 200$ GeV. As m_χ increases to ~ 40 GeV, then, as seen in Fig. 1, our LSP acquires significant shiggs components which can interact with matter via Z° exchange, providing a

larger cross-section. We use the results of Ref. 25), taking $\rho_D^{\text{local}} = 0.3 \text{ GeV/cm}^3$, $\langle v_{\text{rel}} \rangle = 300 \text{ km/s}$ and assuming an idealized detector in which all nuclear recoil energies can be detected. We then estimate the following interaction rates for $v/\bar{v} = 2$, $\lambda = 0.15$ and $m_t = 40 \text{ GeV}$:

$$\begin{aligned} m_\chi &= (20, 30, 40) \text{ GeV} \\ \# \text{events/kg/day} &= (23, 8, 6) \times 10^{-3} \end{aligned} \tag{24}$$

Similar rates are obtained for other values of λ and plausible values of $v/\bar{v}$ in the allowed region of Figs. 2 and 3 where $m_{\frac{1}{2}} \gtrsim 100 \text{ GeV}$, corresponding to $m_\chi \gtrsim 20 \text{ GeV}$. Some proponents²⁵⁾ of dark matter detectors believe that background levels may be as low as $(1 \text{ to } 5) \times 10^{-3} \text{ events/kg/day}$, in which case the rates (24) might be detectable. Clearly the prospects for laboratory detection are brightest if $m_\chi \sim 20 \text{ GeV}$, but one should aim for a detector of mass considerably larger than a kilogram.

We conclude that superstring dark matter may be detectable as well as inevitable.

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	$(ud)_L$	u_L^c	d_L^c	$(\nu e)_L$	e_L^c	D_L	D_L^c	H_0^+	$\bar{H}^-$	ν^c	N
Y	$1/3$	$-4/3$	$2/3$	-1	2	$-2/3$	$2/3$	1	-1	0	0
Y_E	$2/3$	$2/3$	$-1/3$	$-1/3$	$2/3$	$-4/3$	$-1/3$	$-4/3$	$-1/3$	$5/3$	$5/3$
T^3	$\pm 1/2$	0	0	$\pm 1/2$	0	0	0	$\pm 1/2$	$\pm 1/2$	0	0

Table: $U(1)_Y$ quantum numbers (Y), $U(1)_E$ quantum numbers (Y_E), and isospin (T^3) for matter fields arising in the minimal superstring model discussed in this paper.

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FIGURE CAPTIONS

Fig. 1 : The $\tilde{W}^3$, $\tilde{B}$, $\tilde{A}$, $\tilde{S}$, $\tilde{B}_E$ and $\tilde{N}$ content of the LSP as a function of $m_{\frac{1}{2}}$, taking $v/\bar{v} = 2$, $\lambda = 0.15$. The antisymmetric and symmetric shiggs combinations $\tilde{A}$ and $\tilde{S}$ are defined in the text.

Fig. 2 : The mass of the LSP m_χ in GeV (dashed line), its cosmological relic mass density Ω_D in the form $\tilde{\Omega}_D \equiv \Omega_D \times (h_0/0.5)^2$ (solid line), the calculated neutrino fluxes from LSP annihilations in the Sun: $\phi_{\nu_e + \bar{\nu}_e} \times 10^4 \text{cm}^{-2} \text{s}^{-1}$ and $\phi_{\nu_\mu + \bar{\nu}_\mu} \times 10^4 \text{cm}^{-2} \text{s}^{-1}$, and the atmospheric backgrounds $\phi_{\nu_e + \bar{\nu}_e}^{\text{ATM}} \times 10^4 \text{cm}^{-2} \text{s}^{-1}$ and $\phi_{\nu_\mu + \bar{\nu}_\mu}^{\text{ATM}} \times 10^4 \text{cm}^{-2} \text{s}^{-1}$ (dashed-dotted lines), as functions of $m_{\frac{1}{2}}$ for $\lambda = 0.15$, $v/\bar{v} = 2$ and $m_t = 40$ GeV. The background neutrino fluxes were calculated from Ref. 14) assuming $\langle E_\nu \rangle = 0.05 m_\chi$. The down bars on these background flux curves indicate that solar neutrino fluxes at the level of 10% of ϕ^{ATM} should be observable if one exploits directional information in present proton decay detectors. The right-hand scale offers an alternative interpretation of the relic density curves by showing the Hubble parameter $h_0 \equiv H_0/100 \text{ km/s/Mpc}$ predicted by the superstring if one assumes $\Omega_D = 1$. For example, if $m_{\frac{1}{2}} < 240$ GeV, which is required for an acceptable solar neutrino flux if $m_t = 40$ GeV, then $\Omega_D = 1$ requires $h_0 \lesssim 0.45$.

Fig. 3 : Contours of constant $\tilde{\Omega}_D \equiv \Omega_D \times (h_0/0.5)^2$ as a function of λ and $m_{\frac{1}{2}}$. These graphs use $v/\bar{v} = 2$ and $m_t = 40$ GeV, but are quite insensitive to variations in these parameters. The figure shows that $\tilde{\Omega}_D \approx 2^{0 \pm 1}$ over much of the range of parameters $(\lambda, m_{\frac{1}{2}})$, especially when $m_\chi (\approx m_{\frac{1}{2}}/6) \lesssim m_t$ as required for an acceptable solar neutrino flux.

Fig. 1

Fig. 2

Fig. 3