Z_2-Singlino Dark Matter in a Portal-Like Extension of the Minimal Supersymmetric Standard Model

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We propose a Z_2-stabilised *singlino* ($\overline{\chi}$) as a dark matter candidate in extended and R-parity violating versions of the supersymmetric standard model. $\overline{\chi}$ interacts with visible matter via a heavy messenger field S, which results in a supersymmetric version of the Higgs portal interaction. The relic abundance of $\overline{\chi}$ can account for cold dark matter if the messenger mass satisfies $M_S \lesssim 10^4$ GeV. Our model can be implemented in many realistic supersymmetric models such as the NMSSM and nMSSM.

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I. INTRODUCTION

It is well established that visible matter is not sufficient to account for the observed structure of the Universe. This implies the existence of non-baryonic dark matter (DM). Global fits of cosmological parameters can accurately determine the density of DM, albeit indirectly. Measurements of the anisotropy of the cosmic microwave background radiation (CMBR) and of the spatial distribution of galaxies give for the density of DM [1]

$$\Omega_{\rm DM} h^2 = 0.106 \pm 0.008 \,. \tag{1}$$

Identifying the nature of dark matter is a major goal in astroparticle physics. Many particle physics candidates have been proposed in both supersymmetric (SUSY) and non-supersymmetric extensions of the standard model (SM). In either case the stability of DM is ensured by imposing a global symmetry. The simplest global symmetries considered are Z_2 and $U(1)$; see for example [2, 3, 4, 5, 6, 7].

In low energy effective SUSY theories the symmetry is usually R-parity, $(-1)^{(3B+L+2S)}$, which is imposed to conserve baryon (B) and lepton (L) numbers. As a result the stability of proton is ensured. It turns out that R is +1 for all SM fields and -1 for their superpartners. Thus R-parity, which is a Z_2 symmetry, protects the decay of lightest SUSY particle (LSP) to SM particles. As a result the LSP is a good candidate for DM within minimal SUSY standard model (MSSM) and its extensions as long as the conservation of R-parity is ensured.

However, B and L are accidental global symmetries of SM. Thus it is not clear *a priori* that B and L are conserved within the MSSM. If B and L are violated then R-parity is not conserved. Non-conservation of R-parity is one way to generate small neutrino masses [8], which provide solid evidence for phyiscs beyond the SM. Moreover, if R-parity is violated then leptogenesis is possible [9], which explains the small matter anti-matter asymmetry ($O(10^{-10})$) required for successful Big-Bang nucleosynthesis. However, within the MSSM and its extensions there is no well-motivated particle physics candidate for DM in the presence of R-parity violation[1].

In the following we will explore an alternative possibility for the DM candidate in SUSY models, irrespective of whether R-parity is violated or conserved, by introducing a new Z_2 symmetry and additional singlet fields. Singlet extensions of the MSSM are often considered to ensure that the μ parameter is at the electroweak scale [11]. The prime among them are the NMSSM (the Next-to-Minimal SUSY Standard Model) and the nMSSM (the nearly-Minimal SUSY Standard Model). In such models, if R-parity is conserved then the DM candidate can be an R-parity odd singlino [12]. Here we propose an alternative SUSY DM candidate: a Z_2-odd singlino ($\overline{\chi}$) which is stable without requiring R-parity[2].

Beyond considerations of R-parity, Z_2-singlino dark matter is interesting as a SUSY implementation of gauge singlet dark matter. Gauge singlet scalar dark matter interacting via the Higgs portal [14] was first discussed in detail in [4], with a further study presented in [5] and an earlier analysis given in [6]. With the advent of the LHC, Higgs portal couplings to hidden sector particles are of considerable topical interest. The superpotential coupling we will consider here is the natural extension to SUSY of the Higgs portal concept. However, it is necessarily non-renormalisable are a consequence of SUSY, pointing to the existence of further new particles at the TeV scale.

II. MODEL FOR Z_2-SINGLINO DARK MATTER

A. R-parity conserving SUSY

We extend the MSSM by adding a chiral superfield χ and a messenger field S. We also impose an additional Z_2 symmetry under which χ is odd, while all other fields are even. The full

[1] In supergravity (SUGRA) theories, the gravitino can account for dark matter in certain regions of parameter space since its coupling with matter fields is suppressed by the Planck scale [10].

[2] A different Z_2-singlino dark matter model, which is based on a broken $U(1)$ gauge group, was presented in [13].

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superpotential is

$$W = W_{MSSM} + \lambda_1 S\chi\chi + \lambda_2 S H_u H_d + \frac{M_S}{2}S^2 + \frac{M_\chi}{2}\chi^2 , \quad (2)$$

where

$$W_{\rm MSSM} = h^e_{ij} L_i \ell^c_j H_d + h^u_{ij} Q_i u^c H_u + h^d_{ij} Q_i d^c H_d + \mu H_u H_d . \quad (3)$$

In this case the effective superpotential after integrating out S becomes

$$W = W_{\rm MSSM} + \frac{M_\chi}{2}\chi^2 + \frac{f\,\chi^2 H_u H_d}{M_S} , \quad (4)$$

where $f = \lambda_1\lambda_2$. The term with coupling f is the natural generalisation to SUSY of the Higgs portal-type coupling to χ scalars of the form $\chi^\dagger\chi H^\dagger H$ [14]. However, SUSY implies that the Higgs portal interaction is now non-renormalisable. The Lagrangian terms involving the interaction of χ scalars and fermions, to order $1/M_S$, are then

$$\begin{aligned} -\mathcal{L}_\chi \;\supset\; & |M_\chi|^2\chi^\dagger\chi + M_\chi \overline{\chi}.\overline{\chi} + \Bigg[\frac{f\,M_\chi}{M_S}\chi\chi^\dagger H_u H_d \\ & + \frac{f}{M_S}\chi^2 \overline{H}_u.\overline{H}_d + \frac{f}{M_S}\chi H_d \overline{\chi}.\overline{H}_u + \frac{f}{M_S}\chi H_u \overline{\chi}.\overline{H}_d \\ & + \frac{f}{M_S} H_u H_d \overline{\chi}.\overline{\chi} + \text{h.c.}\Bigg] + O(1/M_S^2) , \quad (5) \end{aligned}$$

where χ denotes the scalar and $\overline{\chi}$ the two-component fermion.

B. R-parity violating SUSY

The superpotential involving R-parity non-conserving interactions is:

$$W \supset W_{\not R_p} + \frac{M_\chi}{2}\chi^2 + h_i \frac{\chi^2 L_i H_u}{M_S} , \quad (6)$$

where

$$W_{\not R_p} = \lambda_{ijk} L_i L_j \ell^c_k + \lambda'_{ijk} L_i Q_j d^c_k + \lambda''_{ijk} u^c_i d^c_j d^c_k + \mu'_i L_i H_u \quad (7)$$

is the R-parity non-conserving superpotential in MSSM. The R-parity violating terms in the Lagrangian involving the interaction of χ, to order $1/M_S$, are then given by

$$\begin{aligned} -\mathcal{L}_\chi \;\supset\; & |M_\chi|^2\chi^\dagger\chi + M_\chi \overline{\chi}.\overline{\chi} + \Bigg[\frac{h_i\,M_\chi}{M_S}\chi\chi^\dagger \tilde{L}_i H_u \\ & + \frac{h_i}{M_S}\chi^2 \overline{L}_i.\overline{H}_u + \frac{h_i}{M_S}\chi H_u \overline{\chi}.\overline{L}_i + \frac{h_i}{M_S}\chi \tilde{L}_i \overline{\chi}.\overline{H}_u \\ & + \frac{h_i}{M_S}\tilde{L}_i H_u \overline{\chi}.\overline{\chi} + \text{h.c.}\Bigg] + O(1/M_S^2) , \quad (8) \end{aligned}$$

where $\tilde{L}_i$ is the slepton doublet.

C. Gauge singlet dark matter

Both the scalar and fermion components of the χ superfield are stable due to the Z_2 symmetry and therefore the lightest of these will be a potential DM candidate. In most cases the lightest component will be the fermion, the Z_2-singlino $\overline{\chi}$, since the scalar component will gain additional mass from SUSY breaking. We will therefore focus on the Z_2-singlino as the DM candidate [3]. Its relic abundance will then be determined by the following scattering processes:

$$\begin{aligned} \overline{\chi}\overline{\chi} &\to \text{MSSM fields} \\ \overline{\chi}\overline{\chi} &\to \chi^\dagger\chi \\ \overline{\chi}\chi &\to \text{MSSM fields} . \end{aligned} \quad (9)$$

The latter two processes will be negligible due to Boltzmann suppression if the χ mass is large compared with the $\overline{\chi}$ mass. We will assume this to be the case in the following. Therefore we will only consider the first class of processes when calculating the relic abundance of $\overline{\chi}$.

III. RELIC ABUNDANCE OF Z_2-SINGLINOS

In this section we calculate the relic abundance of $\overline{\chi}$. We first calculate the scattering cross-section times relative velocity for annihilation processes to MSSM final states.

After electroweak symmetry breaking there are five physical Higgs scalar degrees of freedom. In this letter we will consider the physical Higgs scalars to correspond to gauge eigenstates when calculating the cross-sections, with all Goldstone bosons coming from H_d. The physical Higgs scalars are assumed to have a common mass M_H. In addition, we will consider the gaugino and Higgsino gauge eigenstates to correspond to mass eigenstates with a common neutralino mass. A more general analysis will be presented in future work.

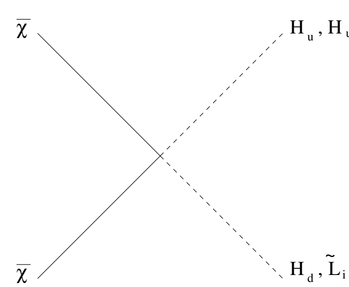

FIG. 1: The four point annihilation of $\overline{\chi}\overline{\chi}$ to Higgs and sleptons in MSSM

In the non-relativistic limit the contribution to the total annihilation cross-section times relative velocity of $\overline{\chi}\overline{\chi}$ annihila-

[3] There may be regions of parameter space where the SUSY mass M_χ is close to the SUSY breaking mass terms, in which case the scalar χ could be the lightest component. We will return to this case in future work.

tion to Higgs and sleptons (Fig.1) is given by:

$$\langle\sigma_1|v_{\rm rel}|\rangle = \frac{1}{4\pi s}\frac{M_{\chi}^2}{M_S^2}\left(1+v_{\rm rel}^2/2\right)\left[f^2\left(1-\frac{2M_H^2}{s}\right) + h_i^2\left(1-\frac{M_{\tilde{L}_i}^2}{s}-\frac{M_H^2}{s}\right)\right]. \qquad (10)$$

The contribution of $\overline{\chi\chi}$ to SM fermions through R-parity con-

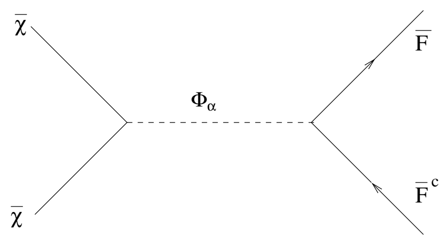

FIG. 2: Mutual annihilation of $\overline{\chi}$ to SM fermions through Higgs and slepton. Here Φ_α=H_u^0, H_d^0, $\tilde{L}_i^0$, $\overline{F}$=$\overline{Q}_i$, $\overline{L}_i$, $\overline{H}_d$ and $\overline{F}^c$=$\overline{u}_j^c$,$\overline{d}_j^c$,$\overline{l}_j^c$.

serving interactions (Fig.2) is given by

$$\langle\sigma_2|v_{\rm rel}|\rangle = \frac{1}{4\pi s}\left(\frac{M_{\chi}^2}{s}\right)\left(1+v_{\rm rel}^2/2\right)\left[\left(\frac{f\langle H_d\rangle}{M_S}\right)^2|h_{ij}^u|^2 \frac{\left(1-\frac{2M_{\tilde{u}}^2}{s}\right)^2}{\left(1-\frac{M_H^2}{s}\right)^2} + \left(\frac{f\langle H_u\rangle}{M_S}\right)^2|h_{ij}^d|^2\frac{\left(1-\frac{2M_{\tilde{d}}^2}{s}\right)^2}{\left(1-\frac{M_H^2}{s}\right)^2} + \left(\frac{f\langle H_u\rangle}{M_S}\right)^2|h_{ij}^e|^2\frac{\left(1-\frac{2M_{\tilde{L}}^2}{s}\right)^2}{\left(1-\frac{M_H^2}{s}\right)^2}\right]. \qquad (11)$$

The contribution of $\overline{\chi\chi}$ to SM fields through R-parity violating interactions (Fig.2) is given by

$$\langle\sigma_3|v_{\rm rel}|\rangle = \frac{1}{4\pi s}\left(\frac{M_{\chi}^2}{s}\right)\left(1+v_{\rm rel}^2/2\right)\left[\left(\frac{h_i\langle H_u\rangle}{M_S}\right)^2|h_{ij}^e|^2\frac{\left(1-\frac{M_{\tilde{L}}^2}{s}-\frac{M_{\overline{H}_d}^2}{s}\right)^2}{\left(1-\frac{M_{\tilde{L}_i}^2}{s}\right)^2} + \left(\frac{h_i\langle H_u\rangle}{M_S}\right)^2|\lambda_{ijk}|^2\frac{\left(1-\frac{2M_{\tilde{L}}^2}{s}\right)^2}{\left(1-\frac{M_{\tilde{L}_i}^2}{s}\right)^2} + \left(\frac{h_i\langle H_u\rangle}{M_S}\right)^2|\lambda'_{ijk}|^2\frac{\left(1-\frac{2M_{\tilde{u}}^2}{s}\right)^2}{\left(1-\frac{M_{\tilde{L}_i}^2}{s}\right)^2}\right] \qquad (12)$$

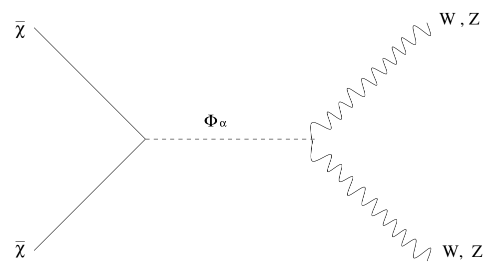

FIG. 3: Mutual annihilation of $\overline{\chi\chi}$ to gauge bosons through Higgs and sleptons. Here $\Phi_\alpha = H_u^0, H_d^0$.

The contribution of $\overline{\chi\chi}$ to W-bosons (Fig.3) is given by

$$\langle\sigma_4|v_{\rm rel}|\rangle = \frac{1}{4\pi}\frac{M_{\chi}^2}{s}\left(1+v_{rel}^2/2\right)\left(2+\frac{(s-2M_W^2)^2}{4M_W^4}\right)\left(1-\frac{2M_W^2}{s}\right)\left[\frac{\left(\frac{f\langle H_d\rangle}{M_S}\right)^2\langle H_u\rangle^2\left(\frac{2M_W^2}{v^2}\right)^2}{(s-M_H^2)^2} + \frac{\left(\frac{f\langle H_u\rangle}{M_S}\right)^2\langle H_d\rangle^2\left(\frac{2M_W^2}{v^2}\right)^2}{(s-M_H^2)^2}\right], \qquad (13)$$

where $v\equiv(\langle H_u\rangle^2+\langle H_d\rangle^2)^{1/2}=246$ GeV, while the contribution of $\overline{\chi\chi}$ to Z-bosons (Fig.3) is given by

$$\langle\sigma_5|v_{\rm rel}|\rangle = \frac{1}{8\pi}\frac{M_{\chi}^2}{s}\left(1+v_{\rm rel}^2/2\right)\left(2+\frac{(s-2M_Z^2)^2}{4M_Z^4}\right)\left(1-\frac{2M_Z^2}{s}\right)\left[\frac{\left(\frac{f^2\langle H_d\rangle}{M_S}\right)^2\langle H_u\rangle^2\left(\frac{2M_Z^2}{v^2}\right)^2}{(s-M_H^2)^2} + \frac{\left(\frac{f^2\langle H_u\rangle}{M_S}\right)^2\langle H_d\rangle^2\left(\frac{2M_Z^2}{v^2}\right)^2}{(s-M_H^2)^2}\right]. \qquad (14)$$

The contribution of $\overline{\chi\chi}$ to sparticles and MSSM Higgs bosons

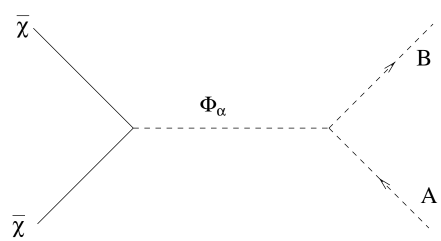

FIG. 4: Mutual annihilation of $\overline{\chi\chi}$ to sparticles and Higgs. Here $\Phi_\alpha = H_u^0, H_d^0, \tilde{L}_i^0$ and A, B stands for the sparticles and Higgs.

(Fig.4) is given by

$$\langle\sigma_6|v_{\rm rel}|\rangle = \frac{1}{4\pi}\frac{M_{\overline{\chi}}^2}{s}\left(1+v_{\rm rel}^2/2\right) \left[\frac{\left(\frac{f\langle H_d\rangle}{M_S}\right)^2}{(s-M_H^2)^2}\sum_{AB}|\mathcal{M}_{AB}|^2\left(1-\frac{M_A^2}{s}-\frac{M_B^2}{s}\right) + \frac{\left(\frac{f\langle H_u\rangle}{M_S}\right)^2}{(s-M_H^2)^2}\sum_{AB}|\mathcal{M}_{AB}|^2\left(1-\frac{M_A^2}{s}-\frac{M_B^2}{s}\right) + \frac{\left(\frac{h_i\langle H_u\rangle}{M_S}\right)^2}{(s-M_{\tilde{L}_i}^2)^2}\sum_{AB}|\mathcal{M}_{AB}|^2\left(1-\frac{M_A^2}{s}-\frac{M_B^2}{s}\right)\right], \quad (15)$$

where $\mathcal{M}_{AB}$ is the mass dimension coupling at the tri-linear scalar vertex. Finally, the contribution of $\overline{\chi}\overline{\chi}$ to gaugino and

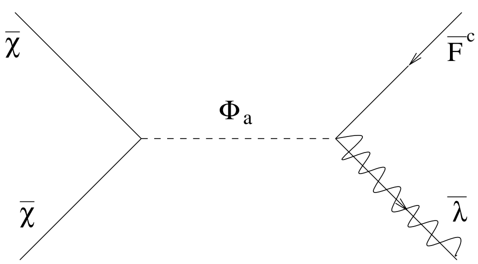

FIG. 5: Annihilation of $\overline{\chi}\overline{\chi}$ to gauginos and fermions. Here $\Phi_\alpha = H_u^0, H_d^0, \tilde{L}_i^0$, $\overline{F} = \overline{H_u}, \overline{H_d}, \overline{L}_i$ and $\overline{\lambda} = \overline{W}, \overline{B}$.

fermion (Fig.5) is given by:

$$\langle\sigma_7|v_{\rm rel}|\rangle = \frac{1}{4\pi s}\left(\frac{M_{\overline{\chi}}^2}{s}\right)\left(1+v_{\rm rel}^2/2\right)\frac{(g^2+g'^2)}{2}\times \left[\left(\frac{f\langle H_d\rangle}{M_S}\right)^2\frac{\left(1-\frac{M_{\overline{H}_u}^2}{s}-\frac{M_{\overline{\lambda}}^2}{s}\right)^2}{\left(1-\frac{M_H^2}{s}\right)^2} + \left(\frac{f\langle H_u\rangle}{M_S}\right)^2\frac{\left(1-\frac{M_{\overline{H}_d}^2}{s}-\frac{M_{\overline{\lambda}}^2}{s}\right)^2}{\left(1-\frac{M_H^2}{s}\right)^2} + \left(\frac{h_i\langle H_u\rangle}{M_S}\right)^2\frac{\left(1-\frac{M_{\overline{L}_i}^2}{s}-\frac{M_{\overline{\lambda}}^2}{s}\right)^2}{\left(1-\frac{M_{\tilde{L}_i}^2}{s}\right)^2}\right]. \quad (16)$$

IV. CONSTRAINTS ON R-PARITY VIOLATING INTERACTIONS AND IMPLICATIONS FOR $\overline{\chi}$ ANNIHILATION

Before estimating the relic abundance of $\overline{\chi}$ let us briefly discuss the constriants on R-parity violating interactions (7) in the MSSM [15]. In MSSM there are three types of trilinear R-parity violating couplings: λ_{ijk}, λ'_{ijk} and λ''_{ijk}. While λ_{ijk} is antisymmetric with respect to i and j, λ''_{ijk} is antisymmetric with respect to j and k. Thus the R-parity violating interactions in general add 45 extra parameters to the MSSM. These couplings are severly constrained by the non-observation of certain physical phenomena. In particular, the product $\lambda'\lambda'' < 10^{-9}$ comes from the stability of proton. Similarly, non-observation of $n-\bar{n}$ oscillations gives the constraint $\lambda'' \leq 10^{-5}$ for $\tilde{m} = 100$ GeV, where $\tilde{m}$ is the SUSY breaking mass. The λ and λ' couplings induce a Majorana mass for three generations of neutrinos. The electron neutrino mass then gives the constraint $\lambda, \lambda' \leq 10^{-3}$ for $\tilde{m} = 100$ GeV. Neutrinoless double beta decay gives the constraint $\lambda' \leq 10^{-4}$. Thus we see that these trilinear couplings are necessarily small in comparison to R-parity conserving couplings in the MSSM. Therefore, the annihilation channels of $\overline{\chi}\overline{\chi}$ through these trilinear R-parity violating couplings are necessarily small in comparison to the R-parity conserving couplings.

There is a bilinear term $\mu'_i L_i H_u$ in the R-parity breaking superpotential. However, one can show that by making a $SU(4)$ rotation $\mu'_i L_i H_u$ can be rotated away [16], leaving only the bilinear term $\mu H_u H_d$ which is R-parity conserving. Therefore the presence of such a bilinear term in the R-parity breaking superpotential does not contribute to any extra annihilations of $\overline{\chi}$.

In what follows we neglect all annihilation channels of $\overline{\chi}\overline{\chi}$ to MSSM fields involving R-parity violating couplings λ, λ' and λ''. However, we note that the new R-parity violating couplings h_i are not necessarily small. When estimating the relic abundance of $\overline{\chi}$ we will consider only those R-parity violating channels involving the couplings h_i.

V. DENSITY OF Z_2-SINGLINO DARK MATTER

The relic abundance of $\overline{\chi}$ can be calculated by solving the Boltzmann equation:

$$\frac{dn_{\overline{\chi}}}{dt} + 3n_{\overline{\chi}}H = -\langle\sigma_{\rm ann}|v_{rel}|\rangle\left(n_{\overline{\chi}}^2 - n_{\overline{\chi}}^{\rm eq2}\right), \quad (17)$$

where $\langle\sigma_{ann}|v_{rel}|\rangle$ is the thermal average of the $\overline{\chi}\overline{\chi}$ annihilation cross-section times relative velocity, with $\sigma_{ann} = \sum_i \sigma_i, i = 1,7$, and $n_{\overline{\chi}}$ is the number density of $\overline{\chi}$. The equilibrium density of non-relativistic $\overline{\chi}$ particles is

$$n_{\overline{\chi}}^{\rm eq} = 2\left[\frac{M_{\overline{\chi}}T}{2\pi}\right]^{3/2} e^{-M_{\overline{\chi}}/T}. \quad (18)$$

With $f = n_{\overline{\chi}}/T^3$, Eq. (17) becomes

$$\frac{df}{dT} = \frac{\langle\sigma_{\rm ann}|v_{rel}|\rangle}{\overline{K}}\left(f^2 - f_{eq}^2\right) , \quad (19)$$

where $f_{eq} = n^{\rm eq}_{\overline{\chi}}/T^3$ and $\overline{K} = \left[4\pi^3 g(T)/45M_{Pl}^2\right]^{1/2}$. The density can then be calculated using the Lee-Weinberg approximation [17]. The freeze-out temperature, T_D, is defined by

$$\frac{df_{eq}}{dT} = \frac{\langle\sigma_{\rm ann}|v_{rel}|\rangle}{\overline{K}} f_{eq}^2 . \quad (20)$$

To obtain the present density Eq. (19) is solved from T_D to the present with $f_{eq} = 0$ on the right-hand side and with $f(T_D) = f_{eq}(T_D)$. The freeze-out temperature can be described by a dimensionless parameter $z_D = M_{\overline{\chi}}/T_D$. Solving Eq. (20) gives for z_D,

$$z_D \equiv \frac{M_{\overline{\chi}}}{T_D} = \ln\left[0.076\frac{1}{g_*^{1/2}}\frac{M_{\overline{\chi}}M_{\rm Pl}\langle\sigma_{\rm ann}|v_{rel}|\rangle}{z_D^{1/2}\left(1-\frac{3}{2z_D}\right)}\right] , \quad (21)$$

where $g_* \equiv g(T_D)$ is the effective number of relativsitic degrees of freedom at T_D. This implies that $z_D \approx 25$. Solving Eq. (19) with $f_{eq} = 0$ on the right-hand side and with $f(T_D) = f_{eq}(T_D)$ then gives the number density at a lower temperature,

$$n_{\overline{\chi}}(T) = \frac{g(T)}{g_*} \times \frac{1.67 g_*^{1/2} T^3 z_D}{M_{\overline{\chi}} M_{Pl}\langle\sigma_{\rm ann}|v_{rel}|\rangle}\frac{\left(1-\frac{3}{2z_D}\right)}{\left(1-\frac{1}{2z_D}\right)} \quad ; \; T \ll T_D , \quad (22)$$

where we have included a correction for the change in the effective number of relativistic degreees of freedom. Therefore the present contribution of $\overline{\chi}$ to the critical density of the universe is

$$\Omega_{\overline{\chi}}h^2 \approx 1.1\times 10^9 {\rm GeV}^{-1}\frac{z_D}{g_*^{1/2}M_{\rm Pl}\langle\sigma_{\rm ann}|v_{rel}|\rangle} , \quad (23)$$

where $z_D \gg 1$ is assumed.

In the following we consider $\Omega_{\overline{\chi}}$ in the limits (i) $s < M_H^2, M_{\tilde{L}}^2$, and (ii) $s > M_H^2$, $M_{\tilde{L}}^2$, where $s \simeq 4M_{\overline{\chi}}^2$ in the non-relativistic limit.

(i) Small $M_{\overline{\chi}}$: $s < M_H^2, M_{\tilde{L}}^2$

To focus on a definite example we set $M_H = M_{\tilde{L}} = 150$ GeV and $\tan\beta \equiv \langle H_u\rangle/\langle H_d\rangle = 1$ in the cross-sections. We assume that the mass of the other sparticles is 100 GeV. Since we assume that $s < M_H^2, M_{\tilde{L}}^2$, in this case only σ_2, σ_3 and σ_7 will contribute to the relic abundance of $\overline{\chi}$. We put $f = h_i = 1$; the results for smaller values can be obtained by rescaling M_S. In this case the allowed region in the plane of M_S versus $M_{\overline{\chi}}$ for $\Omega_{\overline{\chi}}h^2 = 0.106 \pm 0.008$ is shown in Fig. (6). It can be seen that for 15 GeV $\lesssim M_{\overline{\chi}} \lesssim$ 50 GeV, M_S is in the range 1-3 TeV. The behaviour can be understood as follows. In the limit $s < M_H^2, M_{\tilde{L}_i}^2$, the annihilation cross-section $\sigma_{\rm ann} = \sigma_2 + \sigma_3 + \sigma_7$ times relative velocity is of the form:

$$\langle\sigma_{\rm ann}|v_{rel}|\rangle \propto C\frac{M_{\overline{\chi}}^2}{M_S^2} , \quad (24)$$

where C is a dimensionful constant involving the VEV of H_u and H_d. Therefore, smaller values of $M_{\overline{\chi}}$ require small values of M_S in order to keep $\Omega_{\overline{\chi}}h^2$ constant.

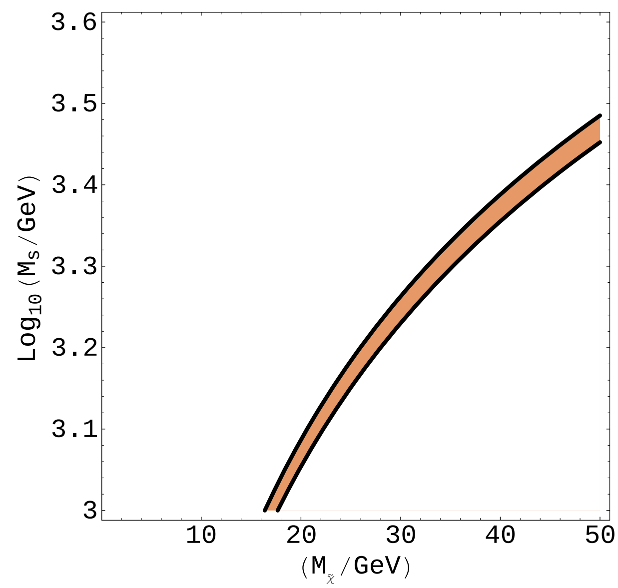

FIG. 6: Contour of $\Omega_{\overline{\chi}}h^2 = 0.106 \pm 0.008$ is shown in the plane of M_S versus $M_{\overline{\chi}}$. We have taken $f = h_i = 1$.

(ii) Large $M_{\overline{\chi}}$: $s > M_H^2, M_{\tilde{L}}^2$

We next consider $s > M_H^2, M_{\tilde{L}}^2$. We show the allowed region in the plane of M_S versus $M_{\overline{\chi}}$, corresponding to $\Omega_{\overline{\chi}}h^2 = 0.106 \pm 0.008$, in Fig. (7). From Fig. (7) it can be seen that for $M_{\overline{\chi}} \gtrsim 200$ GeV, M_S is almost constant at around $10^{3.84}$ GeV $\equiv 6.9$ TeV. This can be understood as follows. In the limit $s > M_H^2, M_{\tilde{L}_i}^2$, the annihilation cross-section $\sigma_{\rm ann} = \sum_i \sigma_i (i = 1-7)$ times relative velocity is of the form:

$$\langle\sigma_{\rm ann}|v_{rel}|\rangle \propto \frac{1}{M_S^2} + C\left(\frac{1}{M_{\overline{\chi}}^2 M_S^2}\right) , \quad (25)$$

where C is a dimensionful constant. For $M_{\overline{\chi}} \gtrsim 200$ GeV, the effective annihilation cross-section is dominated by the first term. As a result we get a constant value $M_S \approx 6.9$ TeV. For $M_{\overline{\chi}} \lesssim 200$ GeV, the second term in the above equation dominates. In this regime, larger M_S is required to keep $\Omega_{\overline{\chi}}h^2$ constant as $M_{\overline{\chi}}$ decreases, with a Higgs pole at $M_{\overline{\chi}} = 75$ GeV allowing much larger values of M_S over a small range of $M_{\overline{\chi}}$.

In general, smaller values of M_S are possible by reducing f and h_i, so the values shown in the figures should be considered as upper bounds on M_S, corresponding to large f and h_i.

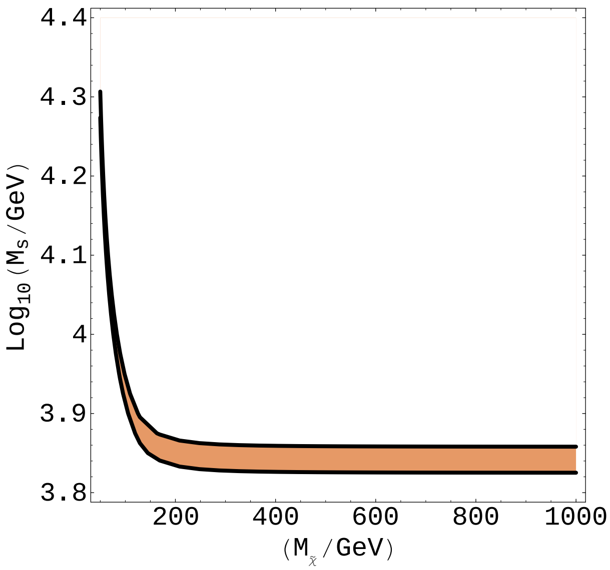

FIG. 7: Allowed region of $\Omega_{\overline{\chi}}h^2 = 0.106 \pm 0.008$ is shown in the plane of M_S versus $M_{\overline{\chi}}$. We have taken $f = h_i = 1$.

VI. CONCLUSIONS AND OUTLOOK

We have discussed the possibility of Z_2-singlino dark matter in extensions of the MSSM. The dark matter communicates with visible matter through a heavy messenger field, S. As a result the interaction is suppressed by the mass scale M_S. For $M_S \lesssim 10^4$ GeV the Z_2-singlino can be cold dark matter for a wide range of mass, 15GeV $\lesssim M_{\overline{\chi}} \lesssim$ 1TeV. (Larger values of M_S are possible near a Higgs pole.) The possibility of dark matter in this case does not rely on the conservation of R-parity. Thus the model is particularly important for the MSSM and its extensions, such as the NMSSM and nMSSM, when R-parity is violated. Non-conservation of R-parity is often considered to give small neutrino masses, as required by the oscillation data, and for leptogenesis, a robust mechanism for the matter anti-matter asymmetry of the Universe.

In the case of non-SUSY gauge singlet scalars interacting via the Higgs portal, direct and indirect detection rates are comparable with conventional weakly interacting dark matter candidates [4, 5]. In the Z_2-singlino case the coupling to the Higgs has an additional suppression factor $\approx v/M_S$, where v is a Higgs expectation value. Therefore we would expect significant detection rates for $M_S \lesssim 1$ TeV. In this case the effective theory based on integrating out the S fields may not be appropriate. We will return to the question of Z_2-singlino detection in future work.

The Z_2 symmetry responsible for dark matter in this model can be a surviving symmetry (a discrete gauge symmetry) of a gauged $U(1)'$ extension of MSSM. Such models are natural in top-down scenarios when $E(6)$ grand unified theory is broken down to the MSSM. A gauge origin of the Z_2 is favoured by arguments which suggest that global symmetries, both continuous and discrete, are broken by non-perturbative gravitational effects [19]. In this case R-parity may be broken while a Z_2 discrete gauge symmetry may account for SUSY dark matter.

We have focused on the case of Z_2-singlino dark matter produced by conventional freeze-out from thermal equilibrium. There is, however, another possibility. In the case of non-SUSY gauge singlet scalar dark matter, when the mass of the scalar is entirely generated by the Higgs expectation value, the correct relic density is produced via decay of thermal background Higgs bosons when the mass of singlet scalars is in the range 1-10 MeV [20]. This is the ideal range [21, 22] for very long-lived dark matter particles to account for the 511 keV line obeserved by INTEGRAL [23]. In the Z_2-singlino model, the singlino mass will be entirely generated by the Higgs expectation value in the limit $M_\chi \to 0$. We will consider the light Z_2-singlino in a forthcoming paper [18].

Although we have considered dark matter particles which are Standard Model singlets, the model can easily be generalised, for example, to a SUSY version of the inert doublet dark matter model [7]. In addition, the messenger mass in the model can be greater than 10^4 GeV, in particular for the case where the Z_2-singlino mass is close to a Higgs pole. This may allow the messengers to be associated with the messenger fields of a gauge mediated SUSY breaking model.

The model we have presented here may be regarded as a SUSY generalisation of the Higgs portal concept. As such, we can expect the model to arise in the low energy effective theory of a wide range of SUSY particle physics models.

Comment: While this paper was in preparation a similar model was presented in [24].

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[1] W. M. Yao et al., Journal of Physics, G33, 1 (2006).
[2] J. McDonald, N. Sahu and U. Sarkar, arXiv:0711.4820 [hep-ph]; S. Nakamura and D. Suematsu, Phys. Rev. D **75**, 055004 (2007).
[3] N. Sahu and U. Sarkar, Phys. Rev. D **76**, 045014 (2007); P. H. Gu and U. Sarkar, Nucl. Phys. B **789**, 245 (2008);

Q. H. Cao, E. Ma and G. Rajasekaran, Phys. Rev. D **76**, 095011 (2007); H. Sung Cheon, S. K. Kang and C. S. Kim, arXiv:0710.2416 [hep-ph].

[4] J. McDonald, Phys. Rev. D **50**, 3637 (1994);

[5] C. P. Burgess, M. Pospelov and T. ter Veldhuis, Nucl. Phys. B **619**, 709 (2001) [arXiv:hep-ph/0011335];

[6] V. Silveira and A. Zee, Phys. Lett. B **161**, 136 (1985).

[7] E. Ma, Phys. Rev. D **73** (2006) 077301; R. Barbieri, L. J. Hall and V. S. Rychkov, Phys. Rev. D **74**, 015007 (2006); L. Lopez Honorez, E. Nezri, J. F. Oliver and M. H. G. Tytgat, JCAP **0702** (2007) 028 [arXiv:hep-ph/0612275]; E. Ma, arXiv:0801.2545 [hep-ph].

[8] R.Hempflingm Nucl. Phys. B478, 3(1996); F.M. Borzumati, Y.Grossman, E. Nardi and Y. Nir, Phys. Lett. B 384, 123 (1996); M. A. Diaz, J.C. Romao and J.W.F. Valle, Nucl. Phys. B524, 23(1998); M.Dress, S. Pakvasa, X. Tata and T. ter Veldhuis, Phys. Rev. D57, 5335 (1998); R. Barbier *et al.*, Phys. Rept. **420** (2005) 1; Y. Grossman and S. Rakshit, Phys. Rev. D **69** (2004) 093002; E. J. Chun, D. W. Jung, S. K. Kang and J. D. Park, Phys. Rev. D **66**, 073003 (2002); A. S. Joshipura, R. D. Vaidya and S. K. Vempati, Nucl. Phys. B **639**, 290 (2002); G. Bhattacharyya, H. V. Klapdor-Kleingrothaus and H. Pas, Phys. Lett. B **463**, 77 (1999).

[9] T. Hambye, E. Ma and U. Sarkar, Nucl. Phys. B **590**, 429 (2000); T. Hambye, E. Ma and U. Sarkar, Phys. Rev. D **62**, 015010 (2000); E. Ma, M. Raidal and U. Sarkar, Phys. Lett. B **460**, 359 (1999);

[10] W. Buchmuller, L. Covi, K. Hamaguchi, A. Ibarra and T. Yanagida, JHEP **0703** (2007) 037 [arXiv:hep-ph/0702184]; and the references there in.

[11] J.E. Kim and H.P. Nilles, Phys. Lett. B **138**, 150 (1984).

[12] C. Hugonie, G. Belanger and A. Pukhov, JCAP **0711**, 009 (2007) [arXiv:0707.0628 [hep-ph]]; G. Belanger, F. Boudjema, C. Hugonie, A. Pukhov and A. Semenov, JCAP **0509**, 001 (2005) [arXiv:hep-ph/0505142]; J. F. Gunion, D. Hooper and B. McElrath, Phys. Rev. D **73**, 015011 (2006) [arXiv:hep-ph/0509024]; D. G. Cerdeno, E. Gabrielli, D. E. Lopez-Fogliani, C. Munoz and A. M. Teixeira, JCAP **0706**, 008 (2007) [arXiv:hep-ph/0701271]; V. Barger, P. Langacker, I. Lewis, M. McCaskey, G. Shaughnessy and B. Yencho, Phys. Rev. D **75**, 115002 (2007) [arXiv:hep-ph/0702036]; F. Ferrer, L. M. Krauss and S. Profumo, Phys. Rev. D **74**, 115007 (2006) [arXiv:hep-ph/0609257].

[13] T. Hur, H. S. Lee and S. Nasri, Phys. Rev. D **77** (2008) 015008 [arXiv:0710.2653 [hep-ph]]; H. S. Lee, arXiv:0802.0506 [hep-ph].

[14] B. Patt and F. Wilczek, arXiv:hep-ph/0605188.

[15] G. Bhattacharyya, Nucl. Phys. Proc. Suppl. **52A**, 83 (1997) [arXiv:hep-ph/9608415].

[16] L. Hall and M. Suzuki, Nucl. Phys. B231, 419 (1984).

[17] B. W. Lee and S. Weinberg, Phys. Rev. Lett. **39** (1977) 165.

[18] J. McDonald and N. Sahu, in preparation.

[19] S. R. Coleman, Nucl. Phys. B **307** (1988) 867.

[20] J. McDonald, Phys. Rev. Lett. **88**, 091304 (2002).

[21] C. Boehm, D. Hooper, J. Silk, M. Casse and J. Paul, Phys. Rev. Lett. **92**, 101301 (2004) [arXiv:astro-ph/0309686].

[22] J. F. Beacom and H. Yuksel, Phys. Rev. Lett. **97**, 071102 (2006) [arXiv:astro-ph/0512411].

[23] P. Jean *et al.*, Astron. Astrophys. **407**, L55 (2003) [arXiv:astro-ph/0309484]; J. Knodlseder *et al.*, Astron. Astrophys. **411**, L457 (2003) [arXiv:astro-ph/0309442].

[24] J. March-Russell, S. M. West, D. Cumberbatch and D. Hooper, arXiv:0801.3440 [hep-ph].