Dark matter candidate with new strong interactions

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We study the possibility that dark matter is a baryon of a new strongly interacting gauge theory, which was introduced in the low energy theory of Cosmological SUSY Breaking (CSB). This particle can fit the observed dark matter density if an appropriate cosmological asymmetry is generated. The same mechanism can also explain the dark/baryonic matter ratio in the universe. The mass of the dark matter particle is in the multiple TeV range, and could be as high as 20 TeV.

I. INTRODUCTION

The most common particle physics models for dark matter involve weakly interacting particles. They can be broadly classified as WIMPS or axions, with the theoretician's favorite WIMP being a neutralino of the Supersymmetric Standard Model (SSM). Within string theory, the physics of both of these candidates is closely connected to SUSY breaking, because string theory axions generally arise from moduli fields, whose mass is related to a superpotential on moduli space.

One of the authors has recently introduced a new model for SUSY breaking, which has no candidate for either WIMP or axion dark matter[1]. The model is based on the principle of Cosmological SUSY Breaking (CSB):

(i) The (positive) cosmological constant (c.c.) is a discrete tunable parameter, governing the number of states in the Hilbert space of quantum gravity in de Sitter (dS) space.

(ii) As the c.c. vanishes, SUSY is restored, with the relation $m_{3/2} \sim \Lambda^{1/4}$ between the gravitino mass and the c.c. A discrete Z_n R symmetry is restored in the same limit, explaining, in low energy terms, the vanishing of the c.c. in the SUSic limiting theory. The limiting theory must have a compact moduli space, in order to guarantee that the dS state of the low energy effective field theory is stable.

(iii) SUSY breaking is spontaneous in the low energy effective theory, but is induced by R breaking terms in the Lagrangian which have no low energy explanation. The coefficients in these terms are tuned to guarantee the CSB scaling relation between $m_{3/2}$ and Λ.

As a consequence of the first requirement, the low energy effective field theory of CSB must contain a goldstino field: a linear supermultiplet which is massless in the SUSic, R symmetric limit. In [1] this was taken to be a chiral superfield G, with R charge 0. If there are no fields of

R charge 2 *mod n* in the low energy theory, then G is naturally massless. R charges were assigned to standard model fields in a way that insured the absence of all baryon and lepton number violating dimension 4 and 5 operators, apart from the term $n^{ij}H_u^2L_iL_j$ (which gives rise to neutrino masses). The generation of this term, and of the texture of Yukawa couplings is imagined to have to do with physics at the unification scale. There is also an ordinary discrete symmetry $\mathcal{F}$, under which G transforms. $\mathcal{F}$ allows the coupling $g_\mu GH_uH_d$ but forbids the conventional μ term. G^a is the lowest order $\mathcal{F}$ invariant monomial in G.

High energy physics supplies us with a term $M_P^2\Lambda^{1/4}f(G/M_P)$ which violates R and implements CSB. The dimensionless coefficients in the function f are tuned to guarantee that the c.c. is indeed Λ. For phenomenological reasons, one must also add terms

$$\int d^4\theta M_1^2 K(g, h_u, h_d, q, \bar{u}, \bar{d}, l, \bar{e}),$$

and,

$$\int d^2\theta Z_A(g^a)W_A^2 + h.c..$$

We have used an unconventional notation where a lower case label s for a chiral superfield S stands for S/M_1. The Kahler potential depends, of course, both on chiral fields and their conjugates. The functions K and Z_A are imagined to emerge from integrating out degrees of freedom at a scale $M_1 \ll M_U \ll M_P$, whose value is determined by RG flow in the limiting $\Lambda = 0$, theory. They can be chosen to satisfy all phenomenological requirements if $M_1 \sim 1$ TeV. It is easy to invent strongly coupled theories $\mathcal{G}$ which could give rise to all the required properties save one. There is no known example of a theory which preserves the R symmetry, and leaves exactly one effective chiral superfield which could play the role of G. We will leave this problem to future work and concentrate on the problem of dark matter.

If the coupling functions Z_A were forced to be logarithms by an accidental $U(1)$ with standard model anoma-

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lies, then the real part of G could be a QCD axion. However, it would have a range of axion couplings ruled out by beam dump experiments. Consequently the model has no axion candidates. The basic setup of CSB contradicts the idea of SUSY neutralino dark matter. The gravitino is the LSP in the CSB scenario, and its longitudinal components are relatively strongly coupled, so the NLSP is not cosmologically stable.

The only plausible dark matter candidate in this scenario is what we will call a G baryon. That is, we assume the strongly interacting G sector has an accidental symmetry, which renders the lightest particle carrying some accidental $U(1)$ quantum number, cosmologically stable. In this paper, we will explore the idea that the dark matter is in fact a baryon of a strongly interacting sector with an RG scale of order M_1. We will see that under a variety of assumptions about the production of this particle, this hypothesis is consistent with conventional cosmology. It has the added virtue of correlating the coincidence between the dynamical scale M_1 and the CSB scale $\sqrt{(\Lambda^{1/4} M_P)}$ to the existence of galaxies. That is to say, we imagine that the limiting model calculates the value of the scale M_1 and the other parameters of e.g. the inflaton field, in such a way that the density of G baryons coincides with what we know about dark matter density from observations. Now consider the model of CSB, with various values of Λ. The only values which will produce a model with galaxies will be those which satisfy Weinberg's bound. At least within a few orders of magnitude, this matches the scale of CSB to M_1 and the dark energy density to the dark matter density (cosmic coincidence).

We will also see that there is a variety of thermal histories for the universe in which G baryons can be dark matter only if there is a CP violating G baryon number asymmetry. We might imagine a model in which G and ordinary baryon asymmetries were produced by the same mechanism, perhaps explaining the dark/baryonic matter ratio of the universe [2,3].

The Hess telescopes [4] have seen a photon signal from the center of the galaxy, which might be consistent with a dark matter candidate of mass 15–18 TeV, if dark matter in the galaxy follows the profile predicted by [5]. It is very hard to find a neutralino model which can produce such a large mass, basically because weak annihilation cross sections decrease with mass. On the other hand, strongly interacting particles have mass independent annihilation cross sections and can easily fit this data.

In the next section, we estimate various cross sections for baryonlike objects, using large N QCD as a paradigm. The G theory must differ from QCD since it preserves chiral symmetry and is supersymmetric. Nonetheless, we hope that these estimates give us a rough guide to the scales involved. We then go on to estimate the mass, cross section and primordial asymmetry for which a G baryon could be dark matter. We consider two scenarios: a standard thermal

relic abundance calculation, and a particular nonthermal production scheme. We find that for reasonable values of parameters, the model can fit the data, and perhaps reproduce the Hess signal. To answer the latter question in more detail, one must perform a detailed estimate of the photon spectrum one would get from annihilation processes involving a strongly interacting dark matter candidate. We are not sure that the model used by the Hess collaboration in order to extract the parameters of a hypothetical dark matter particle from their signal, takes into account the physics of a strongly interacting particle.

We should emphasize that despite our original motivation, our calculations would be applicable to any dark matter candidate with new strong interactions of the right scale. In particular, we note that our model for dark matter is similar to the hypothesis that dark matter is a techni-baryon [6,7].

II. ANNIHILATION CROSS SECTIONS FOR DARK MATTER WITH NEW STRONG INTERACTIONS

The nucleon antinucleon annihilation cross section is usually written in units of the pion Compton wavelength, because this is the range of nuclear forces. In fact, this parametrization is singular in the chiral limit, when the pion becomes a Goldstone boson. It is not correct that the cross section blows up in this limit.

A better estimate is obtained by thinking about chiral soliton models of the nucleon [8]. In such models the nucleon is realized as a classical solution of a large N effective action. The effective Planck constant of this action is of order N, and the scale over which solutions vary is the QCD scale. Although these models use the spontaneously broken chiral symmetry of QCD in an essential way, they give the same order of magnitude results one would expect from general large N considerations. We expect the size of a general large N soliton to be given by such an N independent scale, and large N soliton masses will be of order N.

The soliton-anti-soliton annihilation cross section will be given by its classical size $\sigma \sim \Lambda_G^{-2}$ and will be more or less energy independent in the regime of interest, because the cosmological velocities of these heavy particles will be low. Note that this is *not* s-wave annihilation. The typical orbital angular momentum involved in these collisions is of order $\frac{\sqrt{m_G T}}{\Lambda_G}$, where T is the temperature at which the annihilation takes place. Note also that the thermally averaged cross section $<\sigma v>$, which appears in cosmological Boltzmann equations, will be $O(T/m_G)^{1/2}$. We believe that this is the correct scaling even for ordinary baryons, and that conventional calculations of the relic baryon density in a baryon symmetric universe are not quite correct. However, this does not change the qualitative conclusion of those calculations, namely, that we need a

baryon asymmetry to account for the observed baryon number density of the universe.

We note that the reason that we are interested in large N counting is the combination of the Hess data, and the constraints on Λ_G from supersymmetric phenomenology. The latter prefers a scale $\Lambda_G \sim 1$ TeV, in order to accommodate the bounds on charged superpartner masses, while the former indicates a mass around 15–18 TeV for the dark matter particle. In a large N model, the baryon mass would be $N\alpha\Lambda_G$ with α a number of order 1 ($\alpha \approx 2$ in QCD). Thus, we would want $N\alpha \sim 15 - 18$. These are not unreasonable values. For example, the best of the inadequate models for the G theory, studied in [1] was an $SU(4)$ SUSY gauge theory. For $N = 4$, we require $\alpha \sim 4$, about twice the value in QCD.

We emphasize however that we do not know the details of the model which the Hess collaboration used in quoting 15–18 TeV for their best fit to the dark matter candidate. In particular, for weakly coupled neutral dark matter, the direct photon annihilation signal is suppressed by a power of $(\alpha/\pi)^2$ relative to photons produced from decays of particles with direct coupling to the dark matter. There is no such suppression for strongly interacting neutral composites of charged particles. For example the large N nucleon magnetic moment is order $e(= \sqrt{4\pi\alpha_{em}})N$ in Λ_{QCD} units. Hess has not yet seen the characteristic turnover in their photon signal, which would be expected from dark matter annihilation, and the question of astrophysical explanations for the signal from the galactic center is still controversial. It is perhaps premature to try to fit their spectrum.

However, it is clear that in order to really confront an eventual dark matter signal from Hess data, we need a much better estimate of the photon spectrum produced by a G baryon. In addition, since we find that for most values of the reheat temperature of the universe, we must invoke a G baryon asymmetry to account for the observed dark matter density, the annihilation signal will be proportional to the small density of anti-G baryons. We have not yet done the calculations to determine the range of parameters for which we would expect a significant annihilation signal from the center of the galaxy. In the rest of this paper, we will choose an annihilation cross section of order Λ_G^{-2} and parametrize our results in terms of the G baryon mass $m_G > \Lambda_G$, Λ_G, and an asymmetry.

Our description of the G baryon will utilize the following characteristics of a soliton model: energy independent annihilation cross section much larger than the scale of its Compton wavelength, and thermal production at energies well below its mass. The latter is a well known characteristic [9] of solitons in weakly coupled field theory. Finally, we will parametrize the G baryon mass as $N\alpha\Lambda_G$, with $\Lambda_G \sim 1$ TeV, in order to suggest the large N scaling of soliton masses in strongly coupled gauge theories with large gauge groups.

III. THE RELIC ABUNDANCE OF G BARYONS

We will denote by Ω_G the fraction of the observed density of the universe in G baryons plus antibaryons. To match the observed dark matter abundance, we require $\Omega_G \equiv \frac{\rho_G}{\rho_{cr}} = .24$, using the data from WMAP which specifies $\Omega_m = .29 \pm .07$ and $\Omega_b = .047 \pm .006$. [10] If n_G is the number of G baryons per comoving volume, then this can be written $\Omega = \frac{n_G m_G}{3H_0^2/8\pi G} = \frac{n_G m_G}{1.054 h^2 10^4 \frac{eV}{cm^3}}$.

Writing today's value of the G baryon abundance (the ratio of the number of G baryons per comoving volume and the entropy) $Y_0 \equiv \frac{n_G}{s_0}$, this condition becomes

$$\Omega = .24 = \frac{s_0 Y_0 m_G}{1.054 h^2 10^4 \frac{eV}{cm^3}}$$

Thus we require $Y_0 = \frac{0.44 \ eV}{m_G}$. We will write $m_G = N\alpha$ TeV, treating 1 TeV as the analog of the QCD scale for the G gauge theory, and applying a large N scaling rule for baryon masses. In QCD $N = 3$ and $\alpha \sim 2$. Our point is that the analog of a baryon mass could be quite a bit higher than 1 TeV. For example $N = 6$ and $\alpha \sim 3$ would give us an 18 TeV dark matter candidate, as would be required by the interpretation of Hess data in terms of dark matter annihilation. With this parametrization, the required value of the abundance is $Y_0 \approx \frac{4 \times 10^{-13}}{N\alpha}$.

The relic abundance of G baryons depends on some assumptions about the evolution of the universe at the TeV scale and above. We assume that there was a reheating process which gives rise to a radiation dominated universe at some temperature T_{RH}. This might be due to primordial inflaton decay, or the later decay of some other massive particle which dominates the energy density before it decays. We call the width of the particle Γ_X. If $T_{RH} > 1$ TeV, the G gauge theory is thermalized by X-decay and the post-decay distribution of G baryons is given by the thermal ensemble. Note that this is true even when $m_G \gg T_{RH}$. In this regime of parameters, the G baryon is a thermal relic, and we find that, in the absence of an asymmetry, the relic abundance is too small to explain the observed dark matter density.

For $T_{RH} < 1$ TeV, G baryons are produced nonthermally and we must be a bit more specific about the dynamics. For a weakly coupled X particle, $m_X \gg T_{RH}$ and we can still have X decays into G baryons. Suppose first that $m_X \gg m_G$ so that we can treat the G baryons as just another massless species. If we assume the couplings to G baryons are not suppressed relative to standard model particles we get a branching ratio of order 10^{-2} into G baryons. The decay will be reasonably rapid, so we neglect annihilation processes during the decay period and obtain an initial abundance of

$$Y_0 \sim 10^{-2} \frac{T_{RH}}{m_G}.$$

If the G baryon were massless, this ratio would just be the

branching ratio 10^{-2}. The additional suppression is our estimate of the number of $\mathcal{G}$ baryons per photon that result from the thermalization process.

Throughout the interesting range of parameters, the X particle lifetime is short enough to neglect annihilation in the calculation above. Now we can evolve the resulting $\mathcal{G}$ baryon densities according to the Boltzmann equation driven only by the annihilation of g and $\bar{g}$. As in the discussion of solitonic dark matter abundance in Griest and Kamionkowski [11], the thermally averaged annihilation cross sections have temperature dependence given by: $<\sigma|v|> = \sigma_0(\frac{T_{RH}}{m_G})^{1/2}$. Also, we will assume that $T_{RH} < m_G$. In this case there can be no process in the Boltzmann equation that creates $\mathcal{G}$ baryons because it is not energetically favorable. The Boltzmann equation for the evolution of $\mathcal{G}$ baryons is:

$$\dot{n}_g + 3Hn_g = - <\sigma|v|> n_g^2$$

Letting $Y \equiv n_g/s$ and $x \equiv m_G/T$ we get

$$\frac{dY}{dx} = -\frac{x^{1/2}\sigma_0 s m_{pl}}{1.67 g_*^{1/2} m_G^2} Y^2$$

Since $s = \frac{2\pi^2}{45}g_{*s}T^3$, we can then write:

$$\frac{dY}{dx} = -\frac{kY^2}{x^{5/2}}$$

where $k \equiv \frac{m_G 2\pi^2 \sigma_0 g_{*s} m_{pl}}{1.67 g_*^{1/2} 45}$.

Here we will assume an average $g_* \approx g_{*s} \approx 50$.

Defining an order one parameter β such that $\sigma_0 = \frac{1}{(\beta \text{ TeV})^2}$, $k \approx 4.5 \times 10^{15} N\alpha/\beta^2$.

The solution to this equation is:

$$Y_{\text{final}} = \frac{1}{\frac{1}{Y_i} + \frac{2k}{3}\left(\frac{1}{x_i^{3/2}} - \frac{1}{x_f^{3/2}}\right)}$$

Notice a few properties of this solution. The present day temperature is so low that $\frac{1}{x_f^{3/2}} \approx 0$. Hence either the $\frac{1}{Y_i}$ term or the $\frac{2k}{3}\frac{1}{x_i^{3/2}}$ term dominates, depending on T_{RH}. The $\frac{1}{Y_i}$ term dominates for $T_{RH} < 0.3N\alpha/\beta^2$ MeV. A reheat temperature in this range would be inconsistent with nucleosynthesis, so we can ignore this term. Thus, Y_f is determined by:

$$Y_f = \frac{3x_i^{3/2}}{2k}$$

where $x_i = (m_G/T_{RH})$.

In a general model where we do not fix the mass of the $\mathcal{G}$ baryon or the exact cross section, we can get an upper bound on T_{RH}[1] from our requirement that $Y_0 = \frac{4.4 \times 10^{-13}}{N\alpha}$:

$$T_{RH} > .008\beta^{4/3}N\alpha \text{ TeV}$$

For reheat temperatures below this value, the $\mathcal{G}$ baryons will dominate the universe. Thus we find a small window $1 > \frac{T_{RH}}{\text{TeV}} > 0.008\beta^{4/3}N\alpha$, where nonthermal, symmetric $\mathcal{G}$ baryon production could account for the observed properties of dark matter. In particular, for typical values $N\alpha \sim 10$ and $\beta \sim 1$, we find that this window has a width of about an order of magnitude. However, this range for the reheat temperature does not conform to our prejudice that m_G is substantially larger than Λ_G. We also note that there was no loss of generality in our assumption that $m_X \gg m_G$. If this assumption is not valid, then T_{RH} is quite low, and $\mathcal{G}$ baryons would be overproduced as long as $m_X > m_G$.

For $T_{RH} > 1$ TeV, the thermal relic abundance is too small to account for the observed dark matter, but we can remedy this by postulating an asymmetry. The simplest possibility is that the asymmetry is generated directly in the decay of the X particle, in which case we have the standard result that

$$Y_0 = \epsilon_G \frac{T_{RH}}{m_X},$$

$$\epsilon_G \equiv \sum_f B_f \frac{\Gamma_X(X \to f) - \Gamma_X(X \to \bar{f})}{\Gamma_X}$$

Γ is a decay rate, f and $\bar{f}$ are all possible final states, and B_f is the total $\mathcal{G}$ baryon number of the final state f. ϵ_B is the corresponding asymmetry in ordinary baryon number. In order to match the observed dark matter density and the observed baryon density, we need

$$\frac{\epsilon_G}{\epsilon_B} \approx \frac{1}{2N\alpha} \times 10^{-2},$$

and

$$\epsilon_B \frac{T_{RH}}{m_X} \approx 8.6 \times 10^{-11}.$$

In our model $(\frac{T_{RH}}{m_X} \sim \frac{\sqrt{m_X m_P}}{M})$[2] is bounded from below by the requirements that the X couplings to ordinary matter are at most Planck suppressed, and that the X is massive enough to produce the $\mathcal{G}$ baryon in its decays. Thus $\frac{T_{RH}}{m_X} > (\frac{N\alpha}{2})^{1/2} \times 10^{-3}$.

We see that the ϵ parameters must be very small in order to account for the observed asymmetries. In fact, small baryon number violating branching ratios arise naturally if we assume that X is a slow roll inflaton, I, with a "natural" potential of the form $\mu^4 f(I/m_P)$. A theorem of Nanopoulos and Weinberg [12] tells us that asymmetries can arise only at second order in baryon violating cou-

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[1]When this number is larger than 1 TeV the calculation is not self consistent, because $\mathcal{G}$ baryon production is thermal.

[2]M is the scale of irrelevant couplings of the X particle to the standard model, which are responsible for the decay.

plings. Let us assume the decay of the inflaton is mediated primarily via dimension 5 operators. Then even if the dimension five couplings involve *CP* violation and baryon number violation, we will find that $\epsilon_B \sim (\frac{m_I}{m_P})^2$. We also have the order of magnitude estimate $\frac{T_{RH}}{m_I} \sim (\frac{m_I}{m_P})^{1/2}$, so that

$$Y_B \sim \left(\frac{m_I}{m_P}\right)^{5/2}.$$

This will fit the observed baryon asymmetry if

$$m_I \sim 10^{-4} m_P$$

Note that this gives an inflation scale μ close to the unification scale.

In this context we might attempt to explain the further suppression $\frac{\epsilon_G}{\epsilon_B} \sim 10^{-3}$ by postulating that (perhaps as a consequence of the R symmetry introduced in [1]) the leading contribution to the $\mathcal{G}$-baryon asymmetry comes from the interference of a dimension 5 and dimension 6 coupling of the inflaton, and is suppressed by a further power of $\frac{m_I}{m_P}$. This is off by a factor of 10 but our estimates are so crude that we can consider this a success.

Indeed, we proposed this simple model not because we think it has to be right, but to show that reasonable calculations of both the dark matter and baryon abundances can be obtained for our new form of dark matter.

To summarize, we probably need asymmetric production of $\mathcal{G}$ baryons to make them an acceptable dark matter candidate. We outlined a plausible model of asymmetric production in inflaton decay, which could naturally explain both the baryon asymmetry of the universe and the dark matter density.

IV. CONCLUSIONS

We have shown that a baryonlike state of the new, strongly interacting, $\mathcal{G}$ theory, which was introduced in

[1] to implement Cosmological SUSY breaking, is a promising dark matter candidate. The new strong interaction scale is around 1 TeV and the $\mathcal{G}$ baryon mass is somewhat higher, perhaps as high as the 15–18 TeV needed to fit the Hess data on photons from the center of the galaxy, if the explanation for that data turns out to be dark matter annihilation. We saw that this sort of baryon to interaction scale ratio was natural in the context of large *N* scaling with $N \sim 5-7$. While the $\mathcal{G}$ theory is probably not as simple as $SU(N)$ QCD, there is reason to believe that reasonably large baryon masses are a more general phenomenon.

This sort of dark matter candidate allows one to contemplate a simple explanation of the dark matter to baryon ratio, since the asymmetries in baryon and $\mathcal{G}$ baryon number might have the same physical origin. We need to explain a factor of order 10^3 in these asymmetries, in order to fit the data. We constructed a plausible model in which both asymmetries are generated in inflaton decay. To explain the size of the asymmetries we invoked the R symmetry of [1] and dimensional analysis. There are more scenarios for baryogenesis in the literature than there are authors on this paper, and it is entirely plausible to us that a more elegant mechanism could be found. However, our simple model might work, and it might be the right answer.

Much more work needs to be done to sort out signatures of such hyper-strongly interacting dark matter, as well as to explore a variety of models for the production of baryon and $\mathcal{G}$ baryon asymmetries. In addition, it will be necessary to find out more about the dynamics of the as yet mysterious $\mathcal{G}$ theory, which gives rise to these new particles.

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