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An Inflaton Candidate in Gauge Mediated Supersymmetry Breaking

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Inflation, as currently understood, requires the presence of fields with very flat potentials. Supersymmetric models in which supersymmetry breaking is communicated by supergravity naturally yield such fields, but the scales are typically not suitable for obtaining both sufficient inflation and a suitable fluctuation spectrum. In the context of recent ideas about gauge mediation, there are new candidates for the inflaton. We present a simple model for slow-rollover inflation where the vacuum energy driving inflation is related to the same F-term responsible for the spectrum of supersymmetric particles in gauge mediated supersymmetry breaking models. The inflaton is identified with field responsible for the generation of the μ-term. This opens the possibility of getting some knowledge about the low-energy supersymmetric theory from measurements of the cosmic microwave background radiation. Gravitinos do not pose a cosmological problem, while the moduli problem is ameliorated.

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The existence of an inflationary stage during the evolution of the early Universe is usually invoked to solve the flatness and the horizon problems of the standard big bang cosmology [1]. During inflation the energy density is dominated by the vacuum energy and comoving scales suffer an rapid growth. As a result, any harmful topological defects left as remnants after some Grand Unified phase transition, such as monopoles, are diluted. Quantum fluctuations of the inflaton field imprint a nearly scale invariant spectrum of fluctuations on the background space time metric. These fluctuations may be responsible for the generation of structure formation. However, the level of density and temperature fluctuations observed in the present Universe, $\delta\rho/\rho \sim 10^{-5}$, require the inflaton potential to be extremely flat. For instance, in the chaotic inflationary scenario [2] where the inflaton potential is $V = \lambda\phi^4$ and the condensate sits initially at scales of order of the Planck scale, the dimensionless self-coupling λ must be of order of 10^{-13} to be consistent with observations. The inflaton field must be coupled to other fields in order to ensure the conversion of the vacuum energy into radiation at the end of inflation, but these couplings must be very small otherwise loop corrections to the inflaton potential spoil its flatness. Such flatness is difficult to understand in generic field theories, but is quite common in supersymmetric theories [3,4].

Supersymmetry, then, might be expected to play a fundamental role during inflation [5], but what this role might be will almost certainly depend on the way in which supersymmetry is broken. It is usually assumed that supersymmetry is broken in a hidden sector and then transmitted to observable fields either through gravitational [6] or gauge interactions [7]. The cosmology of this latter class of models has only been partially explored. In this paper we will look for inflaton candidates in this latter class of models, those with gauge-mediated supersymmetry breaking (GMSB) [8].

In this paper we would like to follow the minimal principle that the inflaton field should not be an extra degree of freedom inserted *ad hoc* in some (supersymmetric) theory of particle physics just to drive inflation. We will argue that this is indeed possible in the framework of GMSB models and show that a successful inflation model may be constructed. The energy density driving inflation may correspond to the F-term responsible for the spectrum of the superpartners in the low-energy effective theory. Moreover, we will be able to identify the inflaton field with the same scalar responsible for the generation of the μ-term present in the effective superpotential of the minimal supersymmetric standard model (MSSM), $\delta W \sim \mu H_U H_D$. This raises the possibility of someday connecting a theory which could be tested at accelerators with measurements of the temperature anisotropy in the cosmic microwave background and related measurements of the two-point correlation function! We will see that in such a picture, one can predict the fundamental scale of supersymmetry breaking (which will turn out to be about 10^8 GeV). As it stands, our model is not completely satisfactory. As in virtually all models of inflation based on supersymmetry, a fine tuning is required to obtain a sufficiently small curvature for the potential near the origin. This tuning, at the level of a part in 10^2, must ultimately be given some deeper explanation.

In general, the effective low-energy theory that emerges from GMSB models contains soft SUSY-breaking mass terms for the scalar superpartners which carry information about the scale and nature of the hidden-sector theory. Typical soft breaking terms for sfermions resulting from GMSB models where SUSY is broken at the scale Λ. The magnitude of these terms is given by $\widetilde{m} \sim \frac{\alpha}{4\pi}\Lambda$. This has the interesting consequence that flavour-changing-neutral-current (FCNC) processes are naturally suppressed in agreement with experimental bounds. In the minimal GMSB models, a singlet field X acquires a vacuum expectation value (VEV) for both its scalar and auxiliary components, $\langle X \rangle$ and $\langle F_X \rangle$ respectively and the spectrum for the messangers is rendered non-supersymmetric. Integrating out the messenger sector gives rise to gaugino masses at one loop and scalar masses at two loops, with $\Lambda = F_X/X$. As we shall see, the energy density V_0 driving inflation may be identified with the F-term responsible for the spectrum of supersymmetric particles, $V_0 \sim F_X^2$. This is one of the key results of our paper.

Before launching ourselves into the inflationary aspects of this paper, let us discuss the issue of the generation of the μ-term in the context of GMSB models. Leurer *et al.*, although in a different context, have suggested a μ-term generation which is relevant for us [9]. In addition to the usual MSSM fields, there is another singlet, S. As a consequence of discrete symmetries, the coupling SH_UH_D is forbidden in the superpotential. There are, however, various higher dimension couplings which can drive $\langle S \rangle$. In particular, consider terms in the effective Lagrangian of the form

$$\int d^2\theta \left(\frac{1}{M_p^p} X S^{2+p} + \frac{1}{M_p^m} S^{m+3} + \frac{1}{M_p^n} S^{n+1} H_U H_D \right) + \frac{1}{M_p^2} \int d^4\theta \, X^\dagger X S^\dagger S, \qquad (1)$$

where $M_p \simeq 1.2 \times 10^{19}$ GeV is the Planck mass. This structure can be enforced by discrete symmetries. The first and the fourth terms can contribute to the effective negative curvature terms to the S potential. For example, if $p = m = 2$, and $n = 1$, then the μ-term turns out to be of the order of $\sqrt{F_X}$ (times powers of coupling constants).

The field S, although very weakly coupled to ordinary matter, may play a significant role in cosmology. We will devote the rest of the paper to explore the cosmological implications of such a field and to show that a successfull inflationary scenario may be constructed out of the potential for the field S.

Let us suppose that the field X has acquired a vacuum expectation value (VEV) for both its scalar and auxiliary components and restrict ourselves to the case $p = m = 2$, and $n = 1$. The potential along the real component of the field S reads

$$V(S) \sim \beta^2 F_X^2 - \alpha \frac{F_X^2}{M_p^2} S^2 - \beta \frac{F_X}{M_p^2} S^4 + \beta^2 \frac{X^2}{M_p^4} S^6 - \beta \frac{X}{M_p^4} S^7 + \frac{1}{M_p^4} S^8, \qquad (2)$$

where we have explicitly written the coupling constants. The true vacuum is at $\langle S \rangle^4 \sim \beta F_X M_p^2$, such that $\mu \sim \langle S \rangle^2 / M_p \sim \sqrt{\beta F_X}$. Notice that we have added the constant $\sim F_X^2$ in such a way that the the cosmological constant in the true vacuum is zero, $V(\langle S \rangle) = 0$. As we will see, it is necessary that α be small, of order 10^{-2}. On the other hand, in the framework of supergravity, it is well-known that there are contributions to the mass of particles during inflation of order H from the terms $S^\dagger S$ which must be present in the Kahler potential. These can be cancelled by terms of the form $S^\dagger S X^\dagger X$. This is the fine tuning we have alluded to earlier. It is a generic feature of models based on supersymmetry [10]. One positive feature of these models is that the curvature at the minimum of the potential is automatically far larger than the curvature near the origin. This is important for reheating.

Around $S = 0$ we may simplify the potential as

$$V(S) \simeq V_0 - \frac{m^2}{2} S^2 - \frac{\lambda}{4} S^4, \qquad (3)$$

where $V_0 \sim \beta^2 F_X^2$, $m^2 \sim \alpha \frac{F_X^2}{M_p^2}$ and $\lambda \sim 4\beta \frac{F_X}{M_p^2}$. If the S-field starts sufficiently close to the origin the system may inflate. S rolls very slowly toward $S = \langle S \rangle$ during inflation. Its potential energy dominates the energy density of the Universe driving a nearly constant expansion rate

$$H^2(S) = \frac{8\pi \, V(S)}{3\, M_p^2} \simeq \frac{8\pi \, V_0}{3\, M_p^2}. \qquad (4)$$

In the slow-roll approximation, we may neglect $\ddot{S}$ in the equation of motion so that $\dot{S} \simeq -V'/3H$. This approximation is consistent as long as $V', V'' \ll V$ for $S \sim 0$. Notice that the quadratic and the quartic terms in the expression (4) become comparable for $S_* \sim \sqrt{F_X}$. Since this value is much smaller than $\langle S \rangle$ and all the dynamics giving rise to density perturbations takes place in the vicinity of the origin, we prefer not to neglect the quartic term in the rest of our analysis.

During the slow-roll phase, when S is near to the origin, the cosmic scale factor may grow by $N(S)$ e-foldings:

$$N \simeq \frac{8\pi}{M_p^2} \int_{S_N}^{S_e} \frac{V_0}{-V'(S)} \simeq \frac{2\pi V_0}{m^2 M_p^2} \times \left[1 - 4 \ln \left(\frac{S_N}{S_*} \right) \right], \qquad (5)$$

where S_e denotes the value of the field when inflation ends. Successfull inflation requires $N \simeq 60$. We can express now S_N in terms of S_*: $S_N \simeq S_* \exp\left[1/4 - N(m^2 M_p^2 / 8\pi V_0) \right]$. This expression is valid as long as $S_N < S_*$. The value of the field S_e at which inflation ends corresponds to the moment when the parameter $\epsilon = (M_p^2/4\pi)\,(H'(S)/H(S))^2$ becomes smaller than 1. One can define another fundamental parameter, $\eta = (M_p^2/4\pi)\,(H''(S)/H(S))$. In the slow-roll approximation, these parameters reduce to $\epsilon = (M_p^2/16\pi)(V'/V)^2$ and

$\eta = (M_p^2/8\pi)(V''/V)$ and the slow-roll approximation is valid as long as $\epsilon, \eta \ll 1$. In the model under consideration, however, the slow-roll limit breaks down well before the end of inflation. Using the slow-roll solution for ϵ results in understimating the value S_e at which inflation ends [11], the correct value being S_e of the order of $\langle S \rangle$. Even though the correct value of η becomes large before the end of inflation, the slow-roll approximation is valid for $S \ll \langle S \rangle$, which is the region where observable parameters must be computed [11]. In this region one may compute the power spectrum, which is the Fourier transform of the two-point density autocorrelation function. It has the primordial form $P(k) \propto k^n$. where k is the amplitude of the Fourier wavevector and n denotes the spectral index. Fluctuations arise due to quantum fluctuations in the scalar field S. The measurement of the quadrupole anisotropy in the cosmic microwave background radiation detected by COBE allows us to fix the parameters of the model:

$$\left(\frac{\Delta T}{T}\right) = \sqrt{\frac{32\pi}{45}}\frac{V_0^{3/2}}{V'(S_N)M_p^3} \simeq \frac{\beta^{7/2}}{\alpha^{3/2}}\frac{\sqrt{F_X}}{M_p}\mathrm{e}^{-\frac{1}{4}+N|\eta|}. \tag{6}$$

Imposing $\left(\frac{\Delta T}{T}\right) \simeq 6 \times 10^{-6}$ and fixing $\beta \sim 1$, $\alpha \sim 10^{-2}$ and $\sqrt{F_X} \sim 10^8$ GeV gives $|\eta| \sim 10^{-1}$. It is easy to verify that for such a value of $|\eta|$ the condition $S_N < S_*$ we have used in eq. (7) is indeed satisfied. At a first sight the value $\sqrt{F_X} \sim 10^8$ GeV might seem too large in the framework of GMSB models. However, the spectrum of the superparticles depends on the ratio $\Lambda = F_X/X$. The latter is fixed to be relatively small and in the range $(10 - 10^3)$ TeV. This may be easily accomplished if X acquires a large vacuum expectation value via nonrenormalizable operators. Notice that $\sqrt{F_X} \sim 10^8$ GeV corresponds to a very low value of the Hubble parameter during inflation, $H \sim 1$ MeV. Note also that it implies that there is a small dimensionless number in the coupling of S to the Higgs. This number is comparable to the electron Yukawa coupling.

The COBE satellite measured fluctuations in the cosmic microwave background with a spectral index $n = 1.2 \pm 0.6$ (at 2σ level). In our model the spectral index $n \simeq 1 - 2|\eta| \sim 0.8$ is noticeably smaller than 1 and gives rise to a red spectrum. The amplitude of the gravitational waves produced by quantum fluctuations is far too small to be detected since the variation of the field during inflation is much smaller than M_p [12].

Let us briefly address the issue of the initial condition for the field S. We have assumed that underlying the model are discrete symmetries under which S transforms non-trivially. As a result, $S = 0$ is a special point, and it is natural that S may "sit" at this point initially. This despite the fact that it is very weakly coupled to ordinary matter, and might not be in thermal equilibrium. Many models of slow-rollover inflation require a fine-tuning in the initial value for the field to be successfull and the smaller is the scale of inflation the more severe is the fine-tuning [13]. From Eq. (7) we may infer that, in order to achieve the 60 or so e-foldings of inflation required, the initial value of the scalar field must be less than about $2 \times 10^5 \sqrt{\alpha/\beta}\left(\sqrt{F_X}/10^8\ \mathrm{GeV}\right)$ GeV. As a result, only regions where the initial value of the field is small enough will undergo inflation. These regions have grown exponentially in size and they should occupy most of the physical volume of the Universe. We notice that the small value of the field is not spoiled by quantum fluctuations which are of the order of $H/2\pi \sim 4 \times 10^{-4}\beta\left(\sqrt{F_X}/10^8\ \mathrm{GeV}\right)^2$ GeV. Thermal fluctuations might spoil such a localization since $\langle S^2 \rangle_T^{1/2} \sim T \sim \sqrt{F_X}$. However, the inflaton field is so weakly coupled (its couplings are all suppressed by powers of M_p that it is not in thermal contact with the rest of the Universe [13,14].

After inflation ends, the S-field starts oscillating around the minimum of its potential and the vacuum energy that drives inflation is converted into coherent scalar field oscillations. The Universe undergoes a period of matter domination. Reheating takes place when S decays into light fields, which will eventually thermalize and give rise to a thermal bath of radiation. The reheating temperature T_R is determined by the decay width of the scalar oscillations Γ_S, $T_R \sim 0.1\sqrt{\Gamma_S M_p}$ [15]. At the minimum of its potential, the scalar field has a mass squared $m_S^2 = V''(S) \simeq 10\frac{F_X^{3/2}}{M_p}$. For example, for $\sqrt{F_X} = 10^8$ GeV, $m_S \sim 1$ TeV. The scalar oscillations may decay into light Higgsinos $S \to \psi_{H_u}\psi_{H_d}$ with a rate $\Gamma_S = g^2 m_S/4\pi$ with $g \sim \frac{\mu}{\sqrt{F_X}}\frac{\langle S \rangle}{M_p}$ and we have taken into account that $\mu \ll \sqrt{F_X}$. The resulting reheating temperature is then $T_R \sim 10^{-2}\mu F_X^{1/8} M_p^{-1/4} \simeq 10^2\left(\mu/10^3\ \mathrm{GeV}\right)\left(\sqrt{F_X}/10^8\ \mathrm{GeV}\right)^{1/4}$ MeV, which is large enough to preserve the classical cosmology beginning with the era of nucleosynthesis. It seems difficult, however, to push the reheating temperature above the electroweak scale. Thus, it appears that electroweak baryogenesis is not a viable option for the generation of the baryon asymmetry. On the other hand, the decays of the inflaton themselves might be responsible for the baryon asymmetry [16]. The couplings by which the inflaton decays may contain CP-violation and baryon number violation. In order to produce a baryon asymmetry, we must have baryon number violating operators in the Lagrangian, such as $\delta W \sim \left(\frac{S}{M_p}\right)\bar{U}\bar{D}\bar{D}$. The presence of such operator is compatible with the stability of the proton and the experimental absence of neutron-antineutron oscillations. We can estimate the baryon asymmetry produced by the inflaton decay simply. We assume that the amount of baryon number produced per decay is ε. ε is the product of CP-violating phases δ times the ratio of the baryon number violating decay rate over the total decay rate $\Gamma_B/\Gamma_{tot} \sim 10^{-2}\left(\mu/\sqrt{F_X}\right)^2\left(m_S/\langle S \rangle\right)^2$. The number of massless particles produced per decay is $\sim m_S/T_R$. Plugging in the expected values of the inflaton mass and the reheating temperature for $\sqrt{F_X} \sim 10^8$

GeV, we find a baryon to entropy ratio $B \sim 10^{-8}\delta$, which is compatible with the observed value for $\delta \sim 10^{-2}$.

Another remarkable effect of a late period inflation is that the gravitino problem is solved. As is well known, in GMSB scenarios the gravitino is the lightest supersymmetric particles with mass $m_{3/2} \sim F_X/M_p$ (of order of 1 MeV in our case). If a stable gravitino is thermalized in the early Universe and not diluted by any mechanism, its mass density may exceed the closure limit $\Omega_{3/2} \lesssim 1$. Since the number density of gravitinos is fixed once they are thermalized, the above argument sets a stringent upper bound on the gravitino mass, $m_{3/2} \lesssim 2$ keV (without dilution) [17]. However, gravitinos are efficiently diluted during the inflationary stage driven by the field S and they are not produced in the subsequent stage of reheating. Indeed, light gravitinos (or, better to say, the longitudinal components of them) may be regenerated during reheating either by the decays of sparticles (or particles in the messanger sector) or by scatterings processes. However, the first mechanism requires the reheat temperature to be at least of order of the typical sparticle mass, $\widetilde{m} \sim 100$ GeV, and scattering processes regenerate the light gravitinos only if $T_R \gtrsim 10^2 \left(m_{3/2}/1 \text{ MeV}\right)$ TeV [18]. Since in our scenario the reheat temperature turns out to be very low, $T_R \sim 10^2$ MeV, we may safely conclude that gravitinos were not populating the Universe at the beginning of the radiation era: the gravitino problem is solved by the late stage of inflation.

In models of gauge mediation, if we assume that the underlying theory is a string theory, the cosmological moduli problem is even more severe than in the usual supergravity scenarios [19]. String moduli acquire a mass comparable to $m_{3/2}$ and are stable on cosmological time scales. As a result, their energy density is a severe problem [18]. However, this problem is somewhat ameliorated in our scenario. The period of inflation driven by the field S may take place at a sufficiently late stage of the Universe, $H \lesssim m_{3/2}$, that the number density of string moduli is reduced by a factor $\exp(-3N)$ and by the subsequent entropy production at the reheat stage [20]. It is an attractive feature of the present scenario that this is possible. However, this is not enough, since the minimum of the moduli potential, generically, will be shifted by an amount of order M_p during inflation, as a result of couplings of the moduli

$$\int d^4\theta X^\dagger X f(\mathcal{M}) + S^\dagger S g(\mathcal{M}). \tag{7}$$

So it is probably necessary to find symmetry reasons that the minima coincide to a high degree of accuracy [10].

There has been much discussion of the μ problem in the context of gauge mediation, and it is not clear what is the most satisfactory solution. It is appealing that the solution in [8] also provides an interesting candidate inflaton. We have seen that provided F_X is of suitable size, and provided that certain parameters of order one take suitable values, inflation can occur.

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[1] A. H. Guth, Phys. Rev. D **23**, 347 (1981). For a recent review of inflation and list of references, see A. D. Linde, *Particle Physics and Inflationary Cosmology* (Harwood Academic, New York, 1990).

[2] A.D. Linde, Phys. Lett. **B129** (1993) 177.

[3] J. Ellis, D. Nanopoulos, K. Olive and K. Tamvakis, Phys. Lett. **B118** (1982) 335.

[4] M. Grisaru, W. Siegel and M. Rocek, Nucl. Phys. **B159** (1979) 429.

[5] For a review with extensive references, see D. H. Lyth, Report No. hep-ph/9609431.

[6] For a review, see H. P. Nilles, Phys. Rep. **110** (1984) 1; H.E. Haber and G. L. Kane, Phys. Rep. **117** (1985) 75.

[7] M. Dine, W. Fischler, and M. Srednicki, Nucl. Phys. **B189** (1981) 575; S. Dimopoulos and S. Raby, Nucl. Phys. **B192** (1981) 353; M. Dine and W. Fischler, Phys. Lett. **B110** (1982) 227; M. Dine and M. Srednicki, Nucl. Phys. **B202** (1982) 238; M. Dine and W. Fischler, Nucl. Phys. **B204** (1982) 346; L. Alvarez-Gaumé, M. Claudson, and M. Wise, Nucl. Phys. **B207** (1982) 96; C.R. Nappi and B.A. Ovrut, Phys. Lett. **B113** (1982) 175; S. Dimopoulos and S. Raby, Nucl. Phys. **B219** (1983) 479.

[8] M. Dine, A.E. Nelson, and Y. Shirman, Phys. Rev. **D51** (1995) 1362; M. Dine, A.E. Nelson, Y. Nir, and Y. Shirman, Phys. Rev. **D53** (1996) 2658; M. Dine, hep-ph/9607294.

[9] M. Leurer, Y. Nir and N. Seiberg, Nucl. Phys. **B420** (1994) 468.

[10] Since the non-vanishing vacuum energy density breaks supersymmetry during inflation, supergravity may give an extra contribution $\delta m = cH$ to the inflaton and the moduli masses, see M. Dine, L. Randall and S. Thomas, Phys. Rev. Lett. **75** (1995) 398. We assume here that $c \lesssim \alpha$.

[11] W. Kinney, FERMILAB-PUB-97-042-A preprint, hep-ph/9702427.

[12] D. H. Lyth, Phys. Rev. Lett. **78** (1997) 1861.

[13] L. Knox and M.S. Turner, Phys. Rev. Lett. **70** (1993) 371.

[14] R. Holman, P. Ramond and G.G. Ross, Phys. Lett. **B137** (1984) 343.

[15] For a comprehensive discussion of reheating, see E.W. Kolb and M.S. Turner, *The Early Universe*, Addison-Wesley, Redwood City, CA, 1990.

[16] For a general discussion and references therein, see T. Banks and M. Dine, SCIP-96-31 preprint, hep-ph/9608197.

[17] H. Pagels and J.R. Primack, Phys. Rev. Lett. **48** (1982) 223.

[18] A. de Gouvea, T. Moroi and H. Murayama, LBNL-39753 preprint, hep-ph/9701244.
[19] See T. Banks, M. Berkooz and P.J. Steinhardt, Phys. Rev. **D52** (1995) 705 and references therein.
[20] L. Randall adn S. Thomas, Nucl. Phys. **B449** (1995) 229.