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SU(N,1) INFLATION

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and

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Abstract

We present a simple model for primordial inflation in the context of SU(N,1) no-scale n = 1 supergravity. Because the model at zero temperature very closely resembles global supersymmetry, minima with negative cosmological constants do not exist, and it is easy to have a long inflationary epoch while keeping density perturbations of the right magnitude and satisfying other cosmological constraints. We pay specific attention to satisfying the thermal constraint for inflation, i.e. the existence of a high temperature minimum at the origin.

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Recently, there has been much interest in inflationary models in the context of supergravity theories.[1)-11)] All of these models are examples of primordial inflation[12)] or chaotic inflation[4)] involving SU(5) singlet scalar fields. Models involving just a single scalar field for inflation, the inflaton, are in principle the simplest. The couplings of the inflaton are then determined by the necessary requirements for inflation.[13)] It was noticed,[5)] however, that models involving a single scalar field in supergravity theories utilizing minimal kinetic terms could not simultaneously satisfy the inflationary constraints and the thermal constraint[3)] which requires a high temperature minimum at a point where the zero-temperature potential is flat. In this letter, we will show that they can both be satisfied in models involving non-minimal kinetic terms, in particular the no-scale models of supergravity.[14)-17)]

We begin by reviewing the problem in supergravity models with minimal kinetic terms. These are models in which the Kähler potential is expressed as

$$G = \phi_i^* \phi^i + \ln |F|^2 \tag{1}$$

so that the scalar kinetic terms are

$$\mathcal{L}_{K.T.} = - G^i_j \, (\partial_\mu \phi^j) \, (\partial^\mu \phi_i^*) = - \delta^i_j (\partial_\mu \phi^j) \, (\partial^\mu \phi_i^*) \tag{2}$$

where we have defined $G^i_j \equiv \partial G/\partial\phi^i \partial\phi_i^*$ and $F(\phi^i)$ is the superpotential. The scalar potential is given in terms of G by[18)]

$$V = e^G [G_i (G^{-1})^i_j \, G^j - 3] \tag{3}$$

which for G as in eq. (1) yields the familiar result,

$$V = \exp(\phi^i\phi_i^*)[|\phi_i^*F + F_i|^2 - 3|F|^2] \tag{4}$$

where $F_i \equiv \partial F/\partial\phi^i$.

The breaking of supergravity was most simply accomplished by the inclusion of a term in F such as [19]

$$\mu^2(z + \Delta) \tag{5}$$

where we will refer to z as the Polonyi field. The scale μ was chosen so as to arrange the gauge hierarchy

$$M_W \sim \mu^2/M \tag{6}$$

where $M = Mp/\sqrt{8\pi}$ is the normalized Planck mass. Henceforth we will work in units such that $M = 1$. The Polonyi term is the source of various cosmological problems which we will only comment on below. In order to preserve the gauge hierarchy, one can not assocciate the Polonyi field with the inflaton. This would require[1] $\mu \sim 10^{14}$ GeV as opposed to $\mu \sim 10^{10}$ GeV as given by (6).

Because of the large scale for the inflaton, it has been customary to separate the Polonyi, inflaton, and observable sectors of the theory so that

$$F = F(z) + F(\phi) + F(y_i) \tag{7}$$

where z is the Polonyi field, ϕ is the inflaton and y_i are all other matter and higgs fields. Because we did not want the supersymmetry breaking scale to be influenced by the inflaton, we had the requirement[1]

$$e^{G/2}\, G_\phi = \phi^* F + F_\phi = 0 \tag{8}$$

at the minimum $\langle\phi\rangle = v \sim 0(1)$ after inflation. In order to satisfy (8) and ensure a vanishing cosmological constant at $\langle\phi\rangle = v$ we found[1)]

$$F(v) = F_\phi(v) = 0 \tag{9}$$

The requirements for inflation[13)] set other general constraints on $F(\phi)$. For example, it is clear that at $\phi = 0$ we must have $V(0) > 0$. This implies that

$$F_\phi\,(0) \neq 0. \tag{10}$$

A second requirement is that at the origin, the potential must be flat, i.e. $V_\phi(0) = 0$ or

$$e^{-G}\, V_\phi = G_\phi\,(G_i G^i - 3) + G_{i\phi} G^i + G_\phi = 0 \tag{11}$$

at ϕ=0. We will now see that these two conditions alone are enough to show that a finite temperature minimum is not possible at the origin.

For the case of minimal kinetic terms and in the limit that there are a large number of supermultiplets N of matter fields y^i, the finite temperature corrections to (3) can be written as[8),9)]

$$V_T = (1/12)\ NT^2\, e^G\, [G_i G^i - 2] \tag{12}$$

A finite temperture minimum at $\phi = 0$ will require at least that $(V_T)_\phi = 0$ or

$$12N^{-1}T^{-2}e^{-G}(V_T)_\phi = G_\phi(G_i G^i - 2) + G_{i\phi}G^i + G_\phi = 0 \tag{13}$$

It is clear that the only way to satisfy both (11) and (13) is for $G_\phi = 0$ at $\phi = 0$. This implies however that $F_\phi(0) = 0$ and hence violates the condition (10). The above proof applies only to the large N limit with minimal kinetic terms. The general conditions in which a finite temperature minimum at the origin and a flat potential can exist simultaneously has been examined.[20] In what follows we will look at, in detail, one specific case in which these two conditions are satisfied, i.e. SU(N,1) no-scale supergravity.

Supergravity models based on a non-compact SU(N,1)/SU(N)×U(1) [or SU(1,1) /U(1)] Kähler manifold have attracted interest[14,15,16,17] because they have a flat potential at the tree level. Therefore the minimum of the potential, and hence the scale of supersymmetry breaking, is fixed by radiative corrections. This enables one to determine dynamically the weak interation scale as well as the supersymmetry breaking scale. Thus these "no-scale" models offer a solution to the hierarchy problem. It is also of interest to note that $n \geqslant 4$ extended supergravity theories contain analogous non-compact group structures: SU(1,1)/U(1) in the case of n = 4, SU(5,1)/SU(5)×U(1) in the case of n = 5, etc. Therefore, these models are natural examples of models with non-minimal kinetic terms where one might look for a solution to our inflationary difficulties.

Let us now assume the following form for the Kähler potential[15),16]

$$G = -3 \ln \left(f(z,z^*) - \frac{\phi^*\phi}{3} - \frac{y_i^* y^i}{3}\right) + \ln |F|^2 \tag{14}$$

where we will assume second, third, and fourth derivatives of $f(z,z^*)$ vanish at the physical minimum $z=z_0$ and $|f_z|$ = constant.[16] We assume that ϕ and y are decoupled in the superpotential, as in eq. (7). For the purposes of this discussion we will neglect entirely the matter fields y^i. The kinetic terms that one drives from eq. (2) are not well normalized for the fields z and ϕ. The correctly normalized fields are related to them by the following non-analytic transformations,[14),15)]

$$\xi_R = \frac{1}{\sqrt{6}}\bar{G}, \quad \xi_I = e^{(1/3)\,\bar{G}}\sqrt{3/2}\,(z - z^*)$$

$$\Phi = e^{(1/6)\,\bar{G}}\,\phi, \quad Y^i = e^{(1/6)\,\bar{G}}\,y^i \tag{15}$$

where we have defined $\bar{G} = G - \ln|F|^2$. Assuming $f_{zz^*} = 0$, the scalar potential derived from eq. (14) using eq. (3) is

$$V = e^{(2/3)\,\bar{G}}\,|F_\phi|^2 = e^{\bar{G}}|F_\Phi|^2 \tag{16}$$

This closely resembles the scalar potential for global supersymmetry. We note that in the domain of interest for inflation, $z \sim O(1)$ while $|\phi| << 1$ so that the ϕ dependence in G is negligible.

It is worthwhile at this point to make an important observation. A major problem with the original models of inflation with supergravity[1)-3)] was that with the scalar potential described by (4) it was possible to show[5)] that when the thermal constraint was satisfied, the susy-preserving minimum closest to the origin was always separated by another minimum with negative vacuum energy density. It is now clear that such a disaster is not a general feature in supergravity models, since the form of the scalar potential in eq. (16) is positive semidefinite.

It remains now to show that $F(\phi)$ can be chosen so the $V(\phi,\phi^*)$ will lead to inflation with the correct density perturbation and that the thermal constraints can be satisfied. We will make the illustrative choice[16)] in what follows that $e^{\bar{G}} = 1$.

We begin by exprerssing $F(\phi)$ as a polynominal[1)]

$$F(\phi) = \left[\lambda + \Sigma \frac{\lambda_n}{n+1} \phi^{n+1}\right] m_0^2 \tag{17}$$

so that F_ϕ has the simple form

$$F_\phi = m_0^2 \; \Sigma\lambda_n\phi^n \tag{18}$$

We can assume without loss of generality that the global minimum (or the local minimum closest to the origin) occurs at $\langle\phi\rangle=1$. Then, in order not to add a contribution to the cosmological constant, we require F_ϕ (1) = 0 or

$$\Sigma\ \lambda_n = 0 \tag{19}$$

(Because of the non-minimal nature of G (eq. 14) we no longer have to require F(1) = 0). The condition for flatness at the origin V_ϕ (0) = 0 combined with V(0) > 0 leads to $\lambda_1 = 0$. In addition, a massless inflaton at the origin will lead to the greatest amount of inflation hence we take $\lambda_2 = 0$. The simplest superpotential with inflationary capabilities will then involve only λ_0 and $\lambda_3 = -\lambda_0$ or

$$F(\phi) = m_0^2\ [\phi - (1/4)\ \phi^4] \tag{20}$$

where we have now absorbed λ_0 into m_0. This simple form for F has properties very similar to the model discussed in ref. 6. It differs however, as we will soon see, in allowing for a high temperature minimum at the origin.

The remaining paratemter in $F(\phi)$, m_0, is to be determined by the magnitude of density perturbations.[21] From (16) and (20) we see that the scalar potential is

$$V = m_0^4 \, [1 - (\phi^3 + \phi^{*3}) + (\phi\phi^*)^3] \tag{21}$$

For a potential of the form (21), the density perturbations can be written as[22]

$$\frac{\delta\rho}{\rho} = \frac{1}{(2\pi^3)^{1/2}} \left(\frac{2m_0^4}{H}\right) \ln^2 (Hk^{-1}) \tag{22}$$

where H is the Hubble parameter

$$H^2 = (1/3) \; m_0^4 \tag{23}$$

and k is the wave number of the perturbation. From the condition that $\frac{\delta\rho}{\bar{\rho}} \sim 10^{-4}$ we find that $m_0 \sim 10^{-4}$. If we had chosen $e^G < 1$, m_0 would have been correspondingly larger. The inflationary timescale in this model is just

$$H(\Delta t)_I \simeq 3H^2/V_{\phi\phi}(H) \simeq H/4m_0^4 \sim m_0^{-2} >> 65. \tag{24}$$

Thus there is ample inflaton, with our choice of m_0. We will return later to reheating after inflation in this model. We only comment here that the decay rate for the inflation $\Gamma_\phi \sim M_\phi^3 \sim m_0^6$ leads to a reheat temperature $T_R \sim m_0^3$ and the possibility of a maximum baryon-to-photon ratio,

$$\frac{N_B}{N_\gamma} \sim \frac{T_R}{m_0^2}\ \varepsilon \sim 10^{-4}\ \varepsilon \tag{25}$$

where ε is the magnitude of CP violation during the production of ΔB.

We turn now to the temperature corrections for the theory desribed by (14) and (19). Because of the form of G (i.e. $G^i_j \neq \delta^i_j$) it is no longer possible to use the approximation [8),9)] which led to (12). Instead, we must use the full form for the temperature corrections of $O(T^2)$ given in ref. 8:

$$V_T = (1/24)\ \mathrm{Tr}\ (M_B^2 + 1/2\ M^2{}_F)\ T^2 \tag{26}$$

where

$$\mathrm{Tr}\ M_B^2 = 2V^k_k$$

$$(1/2)\ \mathrm{Tr}\ M_F^2 = e^G[|T^{ij}|^2] - 2e^G \tag{27}$$

$$T^{ij} = G^{ij} + G^i G^j - G^k (G^{-1})^p_k\ G^{ij}_p$$

We note here that these equations can only be used for the fields z, ϕ, and y and not ζ, Φ, and Y.

For the Kähler potential, (14), we can simplify (27) so that[15)]

$$T^{ij} = \frac{F^{*ij}}{F^*} + \frac{2}{3}\left(\overline{G}^i\ \overline{G}^j + \frac{\overline{G}^i F^{*j}}{F^*} + \frac{\overline{G}^j F^{*i}}{F^*}\right) \tag{28}$$

Thus, we arrive at the full temperature dependent potential‡

$$
\begin{aligned}
V_T = \frac{T^2}{24} e^{\bar{G}} \{ & (36|f_z|^4 + \frac{4}{9}(\phi\phi^*)^2 + 8|f_z|^2 \phi\phi^* - 2)\ |F|^2 \\
& + (20|f_z|^2 e^{\bar{G}/3} + 28/9(\phi\phi^*)\ e^{\bar{G}/3} + 4N/3)\ |F_\phi|^2 \\
& + (8/3)\ (\phi\, F_{\phi\phi}\, F^*_\phi + \text{h.c.}) \\
& + (2/3)\ e^{\bar{G}/3}\ (\phi^2\, F_{\phi\phi}\, F^* + \text{h.c.}) \\
& + ([8/9]\phi\phi^* + 8|f_z|^2)\ e^{2\bar{G}/3}\ (\phi\, F_\phi\ F^* + \text{h.c.}) \\
& + 3\ e^{-\bar{G}/3}\ |F_{\phi\phi}|^2 \}
\end{aligned}
\tag{29}
$$

Setting for illustration $e^{\langle\bar{G}\rangle} = 1$ and for $\phi \ll 1$ we have

$$
V_T = \frac{T^2}{24} [(20|f_z|^2 + 4N/3) + (36|f_z|^4 + 8|f_z|^2 + 10/9)\ \phi\phi^* + \ldots]\ m_0^4 \tag{30}
$$

which clearly possesses a minimum at $\phi=0$. We have thus shown that it is possible within the context of supergravity to construct a (zero temperature) potential with a shape desired by inflation with a finite temperature minimum at $\phi = 0$.

Before concluding we wish to comment on two points which were touched on

‡The fermionic contribution to V_T must be computed in terms of correctly normalized fermion fields; see Ref. 15. This requires an additional renormalization by a factor $e^{\bar{G}/3}$ which has been performed in writing the following expression. Because the result depends only on e^G and $|f_z|$, the final form is valid for ϕ as well as Φ.

earlier, 1) the Polonyi field and; 2) reheating and the breaking of SU(5). The Polonyi field has continually been somewhat of an embarrassment cosmologially.[23] The problem was that, left alone, the Polonyi field could not dissipate energy fast enough to settle into its minimum. Inflation made no improvement[24], since the position of the vacuum expectation value of the Polonyi field was dependent on the position of the inflaton and hence had to start rolling again after inflation, giving the original problem[23] back again. Moreover, the low mass of the Polonyi field in the minimal supergravity models with $F(z) = \mu^2 (z+\Delta)$, compared to the large Hubble parameter, resulted in $\langle z^2 \rangle \sim H^4/M_z^2$ due to quantum fluctuations[25] produced during inflation[26]. This ensured that z was far from its minimum, thus thwarting any attempt to solve the Polonyi problem by a devious choice of initial conditions.

In the context of minimal supergravity, two possible solutions[24] have been proposed. Both involve a complicated O'Raifeartaigh sector with three complex Polonyi-esque fields. Rather than comment on these solutions, we will re-examine the situation in the above SU(N,1) model. We observe that by judicious choice of the function $f(z,z^+)$ in the Kähler potential (14), it is possible to fix any desired value of the gravitino mass. At the same time, the mass of one of the two real Polonyi particles if fixed to be $O(m_{3/2})$, while the other remains zero at tree level and is subsquently determined by radiative corrections to be $O(m_W^2/m_p)$. Clearly the former field presents no cosmologial problem ($m_{3/2} >> m_W$ here), while the latter does. It would of course be possible to avoid this remaining difficulty by fixing a large mass for the second Polonyi component. This would indeed solve all cosmological problems, but it would not have the feature [14)-17)] desirable for particle physics of a dynmically determined weak interaction scale. There is an alternative strategy[27] based on SU(N,1) which may also solve both dimensions of the Polonyi problem while retaining this no-scale feature. This is based on models in which the gravitino mass is fixed dynamically by radiative corrections to be $O(m_W^p/m_P^{p-1}$ where $p > 1$. In these models both real components of the complex Polonyi field have enhanced self-couplings of order Λ^4_{QCD} $(z/M_A)^n$: $M_A \approx m_W^{p-1}\ m_P^{2-p}$ which permit a more rapid dissipation of the Polonyi field energy. Determination of the weak interaction scale by radiative corrections is possible[27] in this scenario if $M_A \gtrsim O(10^{12})$GeV. Although $e^{\bar{G}}$ is not fixed in this scenario, as we had assumed before equation (17), it is possible to maintain our assumed form (21) of the inflationary potential by a careful rescaling[14,15] of the supoptential in the Kähler potential (14).

Finally we turn to the questions of reheating and the breaking of SU(5). Earlier on we said that the decay of the inflaton could reheat the Universe up to $T_R \sim M_0^3$. While this is a low value for T_R, it may nevertheless be possible to generate a large enough baryon-to-photon ratio [see eq. (25)]. However, it is easy to conceive of two ways in which this analysis may be modified in our favor. One is to note that fluctuations in our more massive Polonyi field can contribute to reheating. The energy stored in ζ oscillations will be $\sim M_\zeta^2 \langle\zeta^2\rangle$ so that

$$T_R^4 \sim M_\zeta^2 \langle\zeta^2\rangle \tag{31}$$

Because of the large mass of the Polonyi field, even small $\langle\zeta^2\rangle$ will lead to large T_R. For example, if we require only that $T_R \gtrsim 10^4$ GeV, so that $(n_B/n\gamma)_{max} \gtrsim 10^{-6}\ \varepsilon$ for $M_H \sim 10^{10}$GeV, we need only $\langle\zeta^2\rangle > 10^{-56}$ or $(\langle\zeta^2\rangle)^{1/2} > 10^{-20}$ H. We actually get much larger values of $\langle\zeta^2\rangle$. Our previous expression[25] $\langle\zeta^2\rangle \sim H^4/M^2$ only applied in the limit $M \ll H$. In our case[28]

$$\langle\zeta^2\rangle \sim 10^{-2}\, H^2 \tag{32}$$

and therefore we can get ample reheating. If there is any reheating due to ζ_R, it is easy to see that it will involve all matter fields. The ζ_R is present in the exponential in equation (15). Thus for all non-trilinear terms in $F(y^i)$ there will be a direct coupling proportional to $\zeta|F_y|^2$ responsible for ζ decay.

Additionally, we have yet to discuss the connection to SU(5) breaking. It is actually possible to decouple the scale of the slow roll-over, fixed by $\delta\rho/\rho$ to be $H \sim m_0{}^2 \sim 10^{10}$ GeV, from the scale present during the end of inflation.[10] This happens automatically if one couples the inflaton to SU(5) non-singlets in a non-trivial way. Indeed, rather than impose a complete decoupling (7) of the two sectors, it is more natural to consider the most general linked superpotential of the form[10]

$$F(\phi,\sigma) = A(\phi) + B(\phi)\ \sigma^2 + C(\phi)\ \sigma^3 + \ldots \tag{33}$$

where Σ is a 24 of SU(5) and $\sigma^n = \mathrm{Tr}\Sigma^n$. All other observable fields y^i can for our purposs be taken to lie at their minima $y^i \sim 0$. **A priori**, A, B, C, etc. are arbitrary analytic functions of ϕ which represent the (in general non-renomalizable) couplings of ϕ to Σ. Care must be taken to ensure that these couplings do not perturb our hard-won satisfaction (30) of the thermal constraint. Because of these couplings, the curvature of the potential in the Σ direction changes as ϕ rolls, and can become negative for some range of ϕ values, triggering[10] a second order phase transition which breaks SU(5). After the phase transition, the dominant scale of the potential is $m_\Sigma \gg H$. In effect, the inflaton becomes a mixture of ϕ and Σ, and can give adequate reheating through the decays of the Σ - oscillations.

In conclusion, we have demonstrated that one can construct a viable model for inflation in the context of a non-compact SU(N,1) supergravity models. We have shown that it is possible in these models with their non-minimal kinetic terms to arrange a finite temperature minimum at a point where the potential is flat, thus satisfying the thermal constraint. Our inflationary model can be combined with complementary strategies[10,27,28] to solve the Polonyi problem and provide reheating adequate for baryosynthesis.

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ERRATUM:

SU(N,1) Inflation

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Eq. 27 should read:

$$\begin{aligned} \mathrm{Tr}M_B^2 &= 2(G^{-1})^j_k\, V^k_j \\ 1/2\ \mathrm{Tr}M_F^2 &= e^G[(G^{-1})^k_j T^{ij}(G^{-1})^\ell_i\, T_{\ell k} - 2] \\ T^{ij} &= G^{ij} + G^i G^j - G^k(G^{-1})^p_k G^{ij}_p \end{aligned} \tag{27}$$

Eq. 29 should read:

$$V_T = \frac{T^2}{24}\, e^{\bar{G}}[4/3(N+4)e^{-\bar{G}/3}|F_\phi|^2 + 3e^{-2\bar{G}/3}|F_{\phi\phi}|^2 + 2|F|^2] \tag{29}$$

Eq. 30 should read:

$$V_T = \frac{T^2}{24}\,[4/3(N+4) + 2\phi\phi^* + \ldots]m_o^4 \tag{30}$$

Ref. 8 should include:

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All conclusions and results of this paper remain unchanged.