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Origin of structure in supersymmetric quantum cosmology

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Can the imprint of an early supersymmetric quantum cosmological epoch be present in our cosmological observations? Addressing this question is the precise purpose of this paper. Perturbations about a supersymmetric Friedmann-Robertson-Walker FRW model are introduced, in particular by expanding scalar and fermionic fields in adequate harmonics in S^3. The homogeneous and isotropic degrees of freedom are treated exactly, while the others are considered up to quartic order. A set of quantum states is then obtained by employing the supersymmetry and Lorentz constraint equations of this model. Finally, a particular quantum state which has properties typical of the conventional no-boundary (Hartle-Hawking) solution is identified. Its relevance towards a scale-free spectrum of density perturbations is then discussed. [S0556-2821(98)50512-7]

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The objective of this paper is to establish *if* and *how* the inclusion of supersymmetry in a quantum cosmological scenario [1–3] can lead to a scale-free spectrum of density fluctuations. We will construct here a model that describes perturbations about a supersymmetric Friedmann-Robertson-Walker (FRW) minisuperspace with complex scalar fields. By doing so, we will advance some of the ideas previously presented in Refs. [1,2] towards an ''observational context.''

To be more precise, most of the previous research in supersymmetric quantum cosmology (SQC) was aimed at finding quantum states and overcoming consistency problems (see Ref. [1]). No plausible attempt was ever made to find *new* quantum states which would have a physical significance regarding (i) a period of evolution from supersymmetric quantum gravitational physics towards a semiclassical stage, together with (ii) identifying the existence of any quantum state associated to structure formation, (iii) followed by establishing how does conventional quantum cosmology harmonize into this picture, (iv) and hence, determining if a path from supersymmetric quantum cosmology physics down to a classical level can be consistently established.
This paper reports on what is a response regarding (ii) above.

In addition, we are also endorsing quantum supergravity [4,5] as an adequate and more attractive low-energy limit theory. In particular, concerning the study of the very early Universe, instead of using standard gravitational theories with matter fields, but *no* supersymmetry [6].

Our approach is further based in two fundamental elements. First, we subscribe to the idea that the presence of supersymmetry in a quantum universe constitutes an element of the most value. SQC is a framework which is entirely devoted to describe the very early Universe, when quantum gravity effects and supersymmetry are *both* dominant. This contrasts with conventional quantum cosmology, where quantum gravity is present, but *not* supersymmetry. Within this point of view, either (a) supersymmetry has been *entirely* broken, while quantum gravity prevalence continues afterwards, or then, (b) conventional quantum cosmology simply constitutes a *coarse grained* description. In particular, extracted from SQC with some ''averaging'' process, whose physical justification is yet to be established.

Secondly, our research is based in that $N=1$ supergravity [4,5] constitutes a ''square-root'' of gravity [7]. This means that in finding a physical state Ψ, it may be sufficient to solve the Lorentz and supersymmetry constraints of the theory. In fact, the algebra of constraints then implies that Ψ will consequently obey the Hamiltonian constraints. Consequently, this interesting property has been explored in many quantum cosmological cases: see Refs. [1,8–10], where the supersymmetry and Lorentz constraints conducted to simple *first-order* differential equations in the bosonic variables. This advantage contrasts with the situation in nonsupersymmetric quantum cosmology: a *second-order* Wheeler-DeWitt equation has to be solved, employing specific boundary conditions [6,11,12].

The action for our model is then retrieved from the general action of $N=1$ of supergravity with scalar supermultiplets, as represented in Eq. (25.12) of Ref. [5]. Our background supersymmetric minisuperspace is constituted by the gravitational field, which is represented by a tetrad $e_{\mu}^{AA'} = e_{\mu}^{a}\sigma_{a}^{AA'}$ (in two-spinor notation), where [1,2]

$$e_{a\mu}=\begin{pmatrix} N(t) & 0 \\ 0 & a(t)E_{\hat{a}i} \end{pmatrix}, \tag{1}$$

with $\hat{a}$ and i run from 1 to 3, $E_{\hat{a}i}$ is a basis of left-invariant 1-forms on the unit S^3 and $N(t)$, $a(t)$, $\sigma_{a}^{AA'}$ ($A=0,1$) denote, respectively, the lapse function, scale factor, and Infeld-Van der Warden symbols [4,5,2]. In addition, we also have the gravitinos which must have the form (see Refs. [1, 2])

$$\psi^{A}{}_{i}=e^{AA'}{}_{i}\overline{\psi}_{A'}(t), \quad \overline{\psi}^{A'}{}_{i}=e^{AA'}{}_{i}\psi_{A}(t), \tag{2}$$

where $\psi_A, \overline{\psi}_{A'}$ constitute time-dependent spinor fields and $\psi_0^A(t), \overline{\psi}_0^{A'}(t)$ are Lagrange multipliers. A set of time-

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dependent complex scalar fields, $\phi,\overline{\phi}$, and their fermionic superpartners, $\chi_A(t),\overline{\chi}_{A'}(t)$ are also included. Finally, we choose a flat Kähler manifold for the scalar fields.

As far as the perturbations about the background minisuperspace are concerned, we take the scalar fields as

$$\Phi(x_i,t)=\phi(t)+\Sigma_{nlm}f_n^{lm}(t)Q_{lm}^n(x_i), \tag{3}$$

together with its complex conjugate, where the coefficients $f_n^{lm},\overline{f}_n^{lm}$ are functions of the time coordinate t and Q_{lm}^n are standard scalar spherical harmonics on S^3, x_i are coordinates on the three-sphere, and with $n=1,2,3\ldots$, $l=0,\ldots,n-1$, $m=-l,\ldots,l$ [13]. The fermionic superpartners are expanded as [14]

$$\mathbf{X}_A(x_i,t)=\chi_A(t)+a^{-3/2}\Sigma_{mpq}\beta_m^{pq}[s_{mp}(t)\rho_A^{nq}(x_i) + \overline{t}_{mp}(t)\overline{\tau}_A^{mq}(x_i)], \tag{4}$$

together with its Hermitian conjugate, with $m=1,\ldots,\infty$, $p,q=1,\ldots,(m+1)(m+2)$ and, where $\rho_A^{mq},\overline{\rho}_{A'}^{mq},\tau_{A'}^{mq},\overline{\tau}_A^{mq}$ are spinor hyperspherical harmonics on S^3. In addition, the time-dependent coefficients t_{mp},s_{mp} and their Hermitian conjugates are odd elements of a Grassmanian algebra, where the matrix β_{pq}^n satisfy $\beta_{npq}^2=\mathbf{21_n}$.

Inserting now Eqs. (1)–(4) into the general action of $N=1$ supergravity with scalar supermatter [5] and using the properties of the harmonics mentioned in Refs. [13,14], we can obtain (after integration) a reduced action which includes an infinite sum of time-dependent harmonic and Fermi oscillators. The next step is to construct the relevant constraint equations for our model.

In order to write down the supersymmetry constraints, we first need to obtain the Hamiltonian of the theory, which has the form $H=N\mathcal{H}+\psi_0^A S_A+\overline{S}_{A'}\overline{\psi}_0^{A'}+M^{AB}J_{AB}+\overline{M}^{A'B'}\overline{J}_{A'B'}$, where $M^{AB},\overline{M}^{A'B'}$ are additional Lagrange multipliers. $\mathcal{H}$ represents the Hamiltonian constraint, while $S_A,\overline{S}_{A'}$ and $J_{AB},\overline{J}_{A'B'}$ denote, respectively, the supersymmetry and Lorentz constraints. After some suitable redefinitions of the ψ^A and χ^A variables (see Refs. [1,2]), the quantum supersymmetry constraints of the model can be constructed from the coefficients in $\psi_0^A,\overline{\psi}_0^{A'}$ in the Hamiltonian. They take the form $S_A=S_A^{(0)}+S_A^{(\text{pertb.})}$, with

$$S_A^{(0)}=-i\chi_A\frac{\partial}{\partial\phi}-\frac{a\psi_A}{2\sqrt{3}}\frac{\partial}{\partial a}-\sqrt{3}a^2\psi_A-\frac{i}{8}\overline{\phi}\chi^B\chi_B\frac{\partial}{\partial\chi^A} -\frac{i}{4}\overline{\phi}\chi_A\psi^B\frac{\partial}{\partial\psi^B}+\frac{3}{4\sqrt{3}}\psi_A\chi^B\frac{\partial}{\partial\chi^B}+\frac{\psi^B\psi_B}{8\sqrt{3}}\frac{\partial}{\partial\psi^A}, \tag{5}$$

and

$$S_A^{(\text{pertb.})}=\frac{\psi_A}{\sqrt{3}}\Sigma_m\left(s_{mp}\frac{\partial}{\partial s_{mp}}-t_{mp}\frac{\partial}{\partial t_{mp}}\right) +\frac{i}{2}\overline{\phi}\Sigma_m\left(s_m\frac{\partial}{\partial s_m}-t_m\frac{\partial}{\partial t_m}\right)\chi_A -i\chi_A\Sigma_n\frac{\partial}{\partial f_n^{lm}}+2ia^2\Sigma_n\overline{f}_n^{lm}(n+1)\chi_A, \tag{6}$$

together with their Hermitian conjugates, $\overline{S}_A=\overline{S}_A^{(0)}+\overline{S}_A^{(\text{pertb.})}$. $S_A^{(0)},\overline{S}_{A'}^{(0)}$ will denote the supersymmetry constraints of the unperturbed background, while $S_A^{(\text{pertb.})},\overline{S}_A^{(\text{pertb.})}$ correspond to the perturbed sector and have the necessary form to produce the corresponding bosonic Hamiltonian constraint of Ref. [14]. Hereafter, the labels n,l,m and m,p will be denoted simply by n and m, respectively.

At this stage, we introduce a natural ansatz for the wave function of the Universe, which has the form

$$\begin{aligned}\Psi&=A+B\psi^C\psi_C+iC\psi^C\chi_C+D\chi^D\chi_D+E\psi^C\psi_C\chi^D\chi_D\\&=A^{(0)}(a,\phi,\overline{\phi})\Pi_nA^{(n)}(a,\overline{\phi},\phi;f_n\overline{f}_n)\Pi_mA^{(m)}(a,\phi,\overline{\phi},s_m,t_m)+B^{(0)}(a,\phi,\overline{\phi})\Pi_nB^{(n)}(a,\overline{\phi},\phi;f_n\overline{f}_n)\\&\quad\times\Pi_mB^{(m)}(a,\phi,\overline{\phi},s_m,t_m)\psi^C\psi_C+C^{(0)}(a,\phi,\overline{\phi})\Pi_nC^{(n)}(a,\overline{\phi},\phi;f_n\overline{f}_n)\Pi_mC^{(m)}(a,\phi,\overline{\phi},s_m,t_m)\psi^C\chi_C\\&\quad+D^{(0)}(a,\phi,\overline{\phi})\Pi_nD^{(n)}(a,\overline{\phi},\phi;f_n\overline{f}_n)\Pi_mD^{(m)}(a,\phi,\overline{\phi},s_m,t_m)\chi^C\chi_C+E^{(0)}(a,\phi,\overline{\phi})\Pi_nE^{(n)}(a,\overline{\phi},\phi;f_n\overline{f}_n)\\&\quad\times\Pi_mE^{(m)}(a,\phi,\overline{\phi},s_m,t_m)\psi^C\psi_C\chi^D\chi_D,\end{aligned} \tag{7}$$

where each wave functional $A^{(n)},A^{(m)},\ldots,E^{(n)},E^{(m)}$ depends only on the individual perturbation modes f_n or s_m,t_m. Several comments are in order at this point. First, the expression (7) satisfies the Lorentz constraints associated with the unperturbed field variables $\psi_A,\overline{\psi}_{A'},\chi_A$ and $\overline{\chi}_{A'}$: $J_{AB}=\psi_{(A}\overline{\psi}_{B)}-\chi_{(A}\overline{\chi}_{B)}=0$. Second, the perturbation modes of the scalar fields and the fermionic partners do not couple to each other in our approximation and this is also translated in the ansatz (7). In addition, we also follow the approach described in [14], where the coefficients $s_m,t_m,\overline{s}_m,\overline{t}_m$ are taken as invariant under local Lorentz transformation to lowest order in perturbation. We will see ahead that the form of Eqs. (5),(6) together with Eq. (7) will produce consistent solutions. Concerning the coefficients $s_m,t_m,\overline{s}_m,\overline{t}_m$, these are taken as invariant under local Lorentz transformation to lowest order in perturbation (see [14] for a related discussion on this issue). Overall, this approach is fully satisfactory. In fact, we will see in the following, how we can extract a consistent set of solutions from Eqs. (5)–(7).

Now, let us substitute Eq. (7) into the supersymmetry constraint (5),(6) and their Hermitian conjugates. It is important to notice that the terms independent of the perturbation modes have to vanish separately from the ones involving the perturbation modes. Furthermore, the several terms with perturbation modes must also vanish independently—since they do not couple with each other (cf. Refs. [13,14], where this procedure was similarly employed). After having divided $S_A\Psi=0$ and $\overline{S}_A\Psi=0$ by Ψ as given in Eq. (7), we then obtain a set of first-order differential equations. Among them we have:

$$\frac{\partial A}{\partial \phi}-\frac{1}{2}\,\overline{\phi}\Sigma_m\left(s_m\frac{\partial}{\partial s_m}-t_m\frac{\partial}{\partial t_m}\right)A+\Sigma_n\frac{\partial A}{\partial f_n}+2a^2\Sigma_n(n+1)\overline{f}_nA=0, \tag{8}$$

$$\frac{\partial E}{\partial \overline{\phi}}+\frac{1}{2}\,\phi\Sigma_m\left(s_m\frac{\partial}{\partial s_m}-t_m\frac{\partial}{\partial t_m}\right)E+\Sigma_n\frac{\partial E}{\partial \overline{f}_n}-2a^2\Sigma_n(n-1)f_nE=0, \tag{9}$$

$$\frac{a}{\sqrt{3}}\frac{\partial A}{\partial a}+2\sqrt{3}a^2A-\frac{2}{\sqrt{3}}\phi\Sigma_m\left(s_m\frac{\partial}{\partial s_m}-t_m\frac{\partial}{\partial t_m}\right)A=0, \tag{10}$$

$$\frac{a}{\sqrt{3}}\frac{\partial E}{\partial a}-2\sqrt{3}a^2E+\frac{2}{\sqrt{3}}\phi\Sigma_m\left(s_m\frac{\partial}{\partial s_m}-t_m\frac{\partial}{\partial t_m}\right)E=0. \tag{11}$$

Concerning the analysis of a full set of equations, notice that Eqs. (8)–(11) are uncoupled, while the remaining ones constitute coupled partial differential equations. With respect to the former ones, it is straightforward to obtain the following solutions:

$$A^{(0)}=\hat{A}_0^{(0)}\,\frac{e^{-3a^2+\overline{\phi}(2\lambda_1-\Omega_2)-\Omega_2\phi}}{a^{\Omega_1}}, \tag{12}$$

$$A^{(n)}=A_0^{(n)}e^{-\lambda_2\phi+\overline{\phi}(2\lambda_3-\lambda_2)}\times e^{2\lambda_4 f_n-2a^2(n+1)f_n\overline{f}_n-(\Omega_3-\lambda_2)\overline{f}_n+(\Omega_3-\lambda_2)f_n}, \tag{13}$$

$$A^{(m)}=A_0^{(m)}e^{2\lambda_5\phi-C_1\phi\overline{\phi}-\Omega_4\overline{\phi}+\Omega_4\phi}\tilde{A}, \tag{14}$$

$$E^{(0)}=\hat{E}_0^{(0)}\,\frac{e^{3a^2+\phi(2\lambda_6-\Omega_5)-\Omega_5\overline{\phi}}}{a^{\Omega_6}}, \tag{15}$$

$$E^{(n)}=E_0^{(n)}e^{-\lambda_7\overline{\phi}+\phi(2\lambda_8-\lambda_7)}\times e^{2\lambda_9\overline{f}_n+2a^2(n-1)f_n\overline{f}_n-(\Omega_7-\lambda_9)f_n+(\Omega_7-\lambda_9)\overline{f}_n}, \tag{16}$$

$$E^{(m)}=E_0^{(m)}e^{2\lambda_8\overline{\phi}-C_2\phi\overline{\phi}-\Omega_9\phi+\Omega_9\overline{\phi}}\tilde{E}, \tag{17}$$

where $\hat{A}_0^{(0)}=A_0^{(0)}e^{3a_0^2},A_0^{(n)},A_0^{(m)},\hat{E}_0^{(0)}=E_0^{(0)}e^{-3a^2}E_0^{(n)},E_0^{(m)}$ denote integration constants and $\tilde{A}$ and $\tilde{E}\sim s_{mp}$ or t_{mp}. It is important to emphasize the use of $\phi=\phi_1+i\phi_2$ or $\phi=re^{i\theta}$ in the process of integration to decouple the physical degrees of freedom encompassed in $\phi,\overline{\phi}$. Notice as well that $\lambda_1,\lambda_2\ldots$ and C_1,C_2 constitute further integration/separation constants. The quantities $\Omega_1,\Omega_2,\ldots$ represent back reactions of the scalar and fermionic perturbed modes in the homogeneous modes and are assumed to be of a very small value (cf. Refs. [13,14]).

Characteristic features of the no-boundary (Hartle-Hawking) solution are present in the bosonic coefficient E (15)–(17) (see Refs. [13,14,1,2]). This state requires $|\Omega_6|\ll 1$ and the term $e^{-na^2f_n\overline{f}_n}$, $(n\gg 1)$ in Eq. (16) to dominate over the other remaining exponential terms. This is equivalent to assuming that the corresponding separation/integration constants in Eqs. (16),(17) to be very small. It seems that the presence of supersymmetry *selects* a set of solutions, where the no-boundary (Hartle-Hawking) quantum state is mandatory. Finally, it is also important to mention that the states corresponding to $\tilde{E}\sim s_{mp}$ or $\tilde{E}\sim t_{mp}$ mean that these solutions would represent one-particle or one-antiparticle states if we adopt the interpretative framework introduced in Ref. [14].

Concerning the B,C,D coefficients, the corresponding equations lead to integral expressions, similar to the ones in Refs. [1,2]. However, the terms in $f_n,\overline{f}_n$ present in those equations imply that $C^{(n)}=0$ is the only possible solution. Hence, we cannot avoid $C=0$, which is a particularly interesting result.

But do the results hereby presented contribute to our understanding of the very early Universe and if yes, how? Some answers to these enticing questions are advanced in the following.

As both a summary and a point of departure for future research, we presented here an extension of the current repertoire of quantum cosmological models. In particular, incorporating supersymmetry and matter fields expanded in adequate spatial harmonics. As a result, we obtained *new* physical solutions. Among these, we identified one with characteristics typical of the no-boundary proposal [11].

But the most important point brought about in this investigation is that supersymmetric quantum cosmology can constitute an ''observational'' subject—namely, in the sense of making specific predictions for cosmological properties from a quantum description. Thus, this endorses supersymmetry as a mandatory component in any realistic analysis of a quantum universe. Within this context, the answer to those questions above is a *yes*, but where some caution is nevertheless required.

In fact, let us take the bosonic coefficient E [see expressions (15)–(17)], when the term $e^{-na^2f_n\overline{f}_n},(n\gg 1)$ in Eq. (16) is dominant over the other remaining exponential terms. Then, this particular fermionic state implies the following expectation values: $\langle f_n^{(1)}\rangle\sim\langle f_n^{(2)}\rangle\sim n^{-1}a_*^{-2}$ (a_* would be the value of a, when the wavelength of the perturbation modes equal a particular horizon side). Once such conditions have been established, they constitute part of the requirements such that the density perturbations $\delta\rho/\rho$ represent an almost scale-free spectrum of fluctuations, in similarity to what is present in Ref. [4].

In addition, notice that each of the several bosonic amplitudes in Eq. (7) corresponds to a specific quantum scenario for the very early Universe. Supersymmetry seems thus to assign several possible fermionic states with distinct bosonic features, each one leading to different scenarios of evolution.

In particular, we found the Hartle-Hawking quantum state. Since such a state may lead to a satisfactory spectrum of density perturbations [13], our results indicate that supersymmetry within a quantum description of the very early Universe intrinsically contains the relevant seeds for structure formation.

Let us also point out that our model has no potential $V(\phi,\overline{\phi};f_n,\overline{f_n})$ for the homogeneous and inhomogeneous modes. The presence of such potential could induce a transition from a quantum supersymmetric Euclidian phase into an inflationary expansion period. But such potentials will also lead to a mixing in the fermionic sectors of Ψ as present in Eq. (7). In other words, it will imply an additional complex coupling between the equations to solve. Currently, no solutions have yet been found in such a scenario, not even in the corresponding homogeneous sector. Hence, we are dealing with a quantum dominated era of evolution. Eventually, a potential term, e.g., $V(\phi,\overline{\phi})\sim M^2\phi\overline{\phi}$ [1,2,8] will be adequately analyzed within this program [15] and permit us to include a suitable inflationary scenario derived from supergravity (e.g., Ref. [16]). A natural extension will be to expand the tetrad and gravitinos within spherical harmonics. This would further illuminate on the states associated with the inhomogeneous fermionic modes [17].

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