Production of Spin 3/2 Particles from Vacuum Fluctuations

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(Received 9 April 1999)

We study the production of spin 3/2 particles in homogeneous scalar and gravitational backgrounds using the mode-mixing Bogolyubov method. Considering only the helicity $\pm 3/2$ states, we can reduce the problem to a standard Dirac fermion calculation and apply the standard techniques in a straightforward way. As an example we consider a specific supergravity inflationary model and calculate the spectrum of gravitinos created during preheating at the end of inflation.

PACS numbers: 98.80.Cq, 04.62.+v, 04.65.+e

The quantization of fields in the presence of external classical backgrounds leads to interesting phenomena such as the production of particles via the amplification of vacuum fluctuations. This effect has been mainly studied in bosonic models, for example, production of scalars or gravitons in scalar or gravitational backgrounds. In addition, this mechanism for the creation of particles is believed to be responsible for the generation of most of the particles that constitute the present Universe [1], and, in fact, it plays a key role in the modern theories of preheating after inflation. In those models, the energy of the inflaton field is resonantly converted into particles during the period of coherent oscillations after inflation. This parametric resonance phenomenon makes the occupation number of the newly created bosons grow exponentially fast and causes their spectra to be characterized by resonance bands. Recently, the resonant generation of spin 1/2 particles has also been considered in the literature [2]. In these works, it has been shown that the limit on the occupation number imposed by the Pauli exclusion principle is saturated and thus the nonperturbative results deviate considerably from what is expected in a perturbative approach.

In this work we are interested in the creation of spin 3/2 particles through the amplification of vacuum fluctuations. The generation of such particles in the early Universe has traditionally been treated by considering the perturbative decay of other particles [3,4], but not using the nonperturbative approach based on the Bogolyubov transformations technique. Some estimations of the gravitino production during inflation, based on the analogy with Dirac fermions, can be found in [5]. The spin 1/2 case suggests that both approaches can give rise to quite different results. This could be of the utmost importance in the so-called gravitino problem: in supergravity models, the superpartner of the graviton field (gravitino) is described by a spin 3/2 particle. If such particles are created after inflation by some mechanism (particle collision, vacuum fluctuations) they could disrupt primordial nucleosynthesis if they do not decay fast enough, or, if they are stable particles and their masses are high, they could overclose the Universe. In the perturbative approach, these facts impose stringent constraints on both the reheating temperature and the gravitino mass [6].

The calculation of spin 3/2 particle production from vacuum fluctuations is plagued with consistency problems that hamper the quantization of such fields in the presence of external backgrounds. It has been known for a long time [7] that a spin 3/2 particle in scalar, electromagnetic, or gravitational backgrounds can give rise to, apart from algebraic inconsistencies, faster than light propagation modes. This fact completely prevents a consistent quantization in such cases [8]. The only theory in which these problems seem to be absent is supergravity, provided the background fields satisfy the corresponding equations of motion [9]. However, the complicated form of the Rarita-Schwinger equation makes it very difficult to extract explicit results even in simple backgrounds. In this paper we will show that when we consider helicity $\pm 3/2$ states (which dominate the high-energy interactions of gravitinos [3,4]) propagating in arbitrary homogeneous (and isotropic) scalar or gravitational backgrounds, the equations can be reduced to a Dirac-like equation. The quantization can be done along the same lines as for Dirac spinors and therefore the standard Bogolyubov technique [10] can be used to calculate the particle production. We will also show explicitly, within a previously considered supergravity inflationary model, that the expected amplification does take place.

The massive spin 3/2 dynamics in flat space-time is described by the Rarita-Schwinger equation. We will include the scalar field coupling by modifying the mass term (following the notation in [4]):

$$\epsilon^{\mu\nu\rho\sigma}\gamma_5\gamma_\nu\partial_\rho\psi_\sigma + \frac{1}{2}(m_{3/2} - \Phi)[\gamma^\mu, \gamma^\nu]\psi_\nu = 0. \tag{1}$$

As usual in supergravity models we will consider Majorana spinors satisfying $\psi_\mu = C\bar{\psi}_\mu^T$ with $C = i\gamma^2\gamma^0$ the charge conjugation matrix. Contracting this equation with ∂_μ and $\gamma_\lambda \gamma_\mu$ we get

$$- \not{\partial}\Phi \gamma^\nu \psi_\nu + \partial^\mu \Phi \psi_\mu + \frac{1}{2}(m_{3/2} - \Phi)(\not{\partial}\gamma^\nu \psi_\nu - \gamma^\nu \not{\partial}\psi_\nu) = 0 \quad (2)$$

and

$$2i(\partial_\lambda \gamma^\sigma \psi_\sigma - \not{\partial}\psi_\lambda) + (m_{3/2} - \Phi)(\gamma_\lambda \gamma^\nu \psi_\nu + 2\psi_\lambda) = 0\,. \quad (3)$$

Finally contracting this last equation with γ^λ we get

$$i(\not{\partial}\gamma^\sigma \psi_\sigma - \gamma^\lambda \not{\partial}\psi_\lambda) + 3(m_{3/2} - \Phi)\gamma^\mu \psi_\mu = 0\,. \quad (4)$$

When $\Phi = 0$ the three equations (2), (3), and (4) can be written as the Dirac equation plus two constraints, i.e.,

$$(i\not{\partial} - m_{3/2})\psi_\mu = 0\,, \quad (5)$$

$$\gamma^\mu \psi_\mu = 0\,, \quad (6)$$

$$\partial^\mu \psi_\mu = 0\,. \quad (7)$$

The general solution of these equations can be expanded in helicity $l = s/2 + m$ modes:

$$\psi_\mu^{pl}(x) = e^{-ipx} \sum_{s,m} J_{sm} u(\vec{p}, s)\epsilon_\mu(\vec{p}, m)\,, \quad (8)$$

with J_{sm} the Clebsch-Gordan coefficients whose values are $J_{-1-1} = J_{11} = 1$, $J_{-11} = J_{1-1} = 1/\sqrt{3}$, and $J_{-10} = J_{10} = \sqrt{2/3}$. Here $u(\vec{p}, s)$ are spinors with definite helicity $s = \pm 1$ and normalized as $u^\dagger(\vec{p}, r)u(\vec{p}, s) = \delta_{rs}$. If we set $p^\mu = (\omega, p \sin\theta \cos\phi, p \sin\theta \sin\phi, p \cos\theta)$ with $p_\mu p^\mu = m_{3/2}^2$ and $p = |\vec{p}|$, then the three spin 1 polarization vectors are given by

$$\epsilon_\mu(\vec{p}, 1) = \frac{1}{\sqrt{2}}(0, \cos\theta \cos\phi - i \sin\phi, \cos\theta \sin\phi + i \cos\phi, -\sin\theta)\,, \quad (9)$$

$$\epsilon_\mu(\vec{p}, 0) = \frac{1}{m_{3/2}}(p, -\omega \sin\theta \cos\phi\,, - \omega \sin\theta \sin\phi, -\omega \cos\theta)\,, \quad (10)$$

$$\epsilon_\mu(\vec{p}, -1) = -\frac{1}{\sqrt{2}}(0, \cos\theta \cos\phi + i \sin\phi, \cos\theta \sin\phi - i \cos\phi, -\sin\theta)\,. \quad (11)$$

The normalization is $\epsilon_\mu^*(\vec{p}, m)\epsilon^\mu(\vec{p}, n) = \delta_{mn}$, $p^\mu \epsilon_\mu(\vec{p}, m) = p^\mu \epsilon_\mu^*(\vec{p}, m) = 0$. The corresponding quantization details can be found elsewhere [4].

Now we turn to the $\Phi \neq 0$ case. The expression in (8) is no longer a solution of the equations of motion. Let us now concentrate on homogeneous scalar fields, dependent only on the time coordinate $\Phi(t)$. We look for general homogeneous solutions of the Rarita-Schwinger equation of the form

$$\psi_\mu^{pl}(x) = e^{i\vec{p}\cdot\vec{x}} f^{pl}(t) \sum_{s,m} J_{sm} u(\vec{p}, s)\epsilon_\mu(\vec{p}, m)\,. \quad (12)$$

These fields satisfy the condition $\gamma^\mu \psi_\mu = 0$, since they differ from (8) in just a scalar factor. Now if we restrict ourselves to the helicity $l = \pm 3/2$ states, they satisfy $\psi_0^{p \pm 3/2} = 0$ and, since the spatial derivatives of the scalar field vanish ($\partial_i \Phi = 0$), then (2) and (4) are automatically satisfied provided $\partial^i \psi_i = 0$. From (12) this last condition is equivalent to $p^i \psi_i = 0$ which holds from the condition $p^\mu \epsilon_\mu(\vec{p}, m) = 0$. Accordingly, for helicity $\pm 3/2$ states propagating in an homogeneous scalar background, the Rarita-Schwinger equation reduces again to the Dirac form

$$[i\not{\partial} - m_{3/2} + \Phi(t)]\psi_\mu^{\pm 3/2} = 0\,. \quad (13)$$

As far as these modes satisfy a Dirac-like equation, it appears that all the difficulties in the quantization would concern just the helicity $\pm 1/2$ modes in this case. In fact, the above ansatz (12) is not a solution for the helicity $\pm 1/2$ modes even for homogeneous backgrounds.

Let us include the effect of curved space-time. We will concentrate on spatially flat Friedmann-Robertson-Walker (FRW) metrics, and we will introduce it by *minimal* coupling as done in supergravity, i.e., $D_\rho \psi_\sigma = (\partial_\rho + \frac{i}{2}\Omega_\rho^{ab}\Sigma_{ab})\psi_\sigma$ with Ω_ρ^{ab} the spin-connection coefficients and $\Sigma_{ab} = \frac{i}{4}[\gamma_a, \gamma_b]$. The $\epsilon^{\mu\nu\rho\sigma}$ removes the Christoffel symbols contribution in the covariant derivative. We will continue considering $\Phi(t)$ to be a function of time alone. We will consider only the linearized equation in $1/M$ (where $M_P^2 = 8\pi M^2$) for supergravity [11], i.e., we will ignore the torsion contribution to the spin connection which is of $\mathcal{O}(M^{-2})$. In this case the equations of motion for the gravitino read

$$\epsilon^{\mu\nu\rho\sigma}\gamma_5 \gamma_\nu D_\rho \psi_\sigma + \frac{1}{2}(m_{3/2} - \Phi)[\gamma^\mu, \gamma^\nu]\psi_\nu = 0\,. \quad (14)$$

Contracting with D_μ, taking into account that $D_\mu \gamma_\nu = 0$ and $[D_\mu, D_\rho] = -\frac{i}{2}R_{\mu\rho}^{ab}\Sigma_{ab}$ (the vector part cancels because of the $\epsilon^{\mu\nu\rho\sigma}$ term), we get

$$-\frac{i}{4}\epsilon^{\mu\nu\rho\sigma}\gamma_5 \gamma_\nu R_{\mu\rho}^{ab}\Sigma_{ab}\psi_\sigma - (\not{D}\Phi)\gamma^\nu \psi_\nu + (D^\mu \Phi)\psi_\mu + \frac{1}{2}(m_{3/2} - \Phi)(\not{D}\gamma^\nu \psi_\nu - \gamma^\nu \not{D}\psi_\nu) = 0\,. \quad (15)$$

Following the same steps as in flat space-time we obtain (3) and (4) but replacing ordinary derivatives by covariant ones. For FRW metrics and helicity $\pm 3/2$ states, i.e., $\psi_0 = 0$, it is possible to show that due to the form of the Riemann tensor the first term in (15) is proportional to $\gamma^\mu \psi_\mu$ and accordingly we get

$$(i\not{D} - m_{3/2} + \Phi)\psi_\mu = 0\,, \quad (16)$$

$$\gamma^\mu \psi_\mu = 0\,, \quad (17)$$

$$D^\mu \psi_\mu = 0\,. \quad (18)$$

Here again we can use the standard formulas for particle production obtained for the spin $1/2$ cases to study the creation of helicity $\pm 3/2$ states in a FRW background. With that purpose we have to reduce Eq. (16) to a second order equation. Let us first write the equation in conformal time defined as $dt = a(\eta)d\eta$:

$$\left(ia^{-1}\delta_a^\mu \gamma^a \partial_\mu - m_{3/2} + \Phi + i\,\frac{3}{2}\,\frac{\dot a}{a^2}\,\gamma^0\right)\psi_\mu = 0\,, \tag{19}$$

where $\dot a = da/d\eta$. We will adopt the following ansatz for the helicity $l = \pm 3/2$ solutions:

$$\psi_\mu^{pl}(x) = a^{-3/2}(\eta) e^{i\vec p\cdot\vec x} U_\mu^{\vec p l}(\eta)\,, \tag{20}$$

with

$$U_\mu^{\vec p l}(\eta) = \frac{1}{\sqrt{\omega + m_{3/2}^0}}\,(i\gamma^0\partial_0 - \vec p\cdot\vec\gamma + a(\eta)[m_{3/2} - \Phi(\eta)]) \times f_{pl}(\eta) u(\vec p, s)\delta_\mu^a \epsilon_a(\vec p, m)\,, \tag{21}$$

and the normalization $U_\mu^{\vec p l\dagger}(0) U_{\vec p l}^\mu(0) = 2\omega$ and $m_{3/2}^0 = a(0)m_{3/2}$. One can check that this ansatz automatically satisfies (17) and (18). An appropriate form for the spinor $u(\vec p, s)$ and polarization vectors $\epsilon_a(\vec p, m)$ can be obtained if we choose the Dirac representation for the gamma matrices and we take (without loss of generality) the z axis to be along the $\vec p$ direction. In this case $u(\vec p, 1)^T = (1,0,0,0)$, $u(\vec p, -1)^T = (0,1,0,0)$, $\epsilon_a(\vec p, 1) = \frac{1}{\sqrt2}(0,1,i,0)$, and $\epsilon_a(\vec p, -1) = \frac{1}{\sqrt2}(0,1,-i,0)$. With this choice, $u(\vec p, \pm1)$ are eigenstates of γ^0 with eigenvalues $+1$. Then Eq. (19) reduces to the well-known form

$$\left(\frac{d^2}{d\eta^2} + p^2 - i\,\frac{d}{d\eta}\,(a(\eta)[m_{3/2} - \Phi(\eta)]) + a^2(\eta)\,(m_{3/2} - \Phi(\eta))^2\right) f_{pl}(\eta) = 0\,. \tag{22}$$

In order to quantize the modes we will expand an arbitrary solution with helicity $l = \pm 3/2$ as

$$\psi_\mu^l(x) = \int \frac{d^3p}{(2\pi)^3 2\omega}\, a^{-3/2}(\eta) \times [e^{i\vec p\cdot\vec x} U_\mu^{\vec p l}(\eta) a_{\vec p l}) + e^{-i\vec p\cdot\vec x} U_\mu^{\vec p lC}(\eta) a_{\vec p l}^\dagger]\,, \tag{23}$$

where the creation and annhilation operators satisfy the anticommutation relations $\{a_{\vec p l}, a^\dagger_{\vec p' l'}\} = (2\pi)^3 2\omega\,\delta_{ll'}\delta(\vec p - \vec p')$.

In order to see how this works in practice, we will consider a specific supergravity inflationary model (see [12,13]), in which the inflaton field is taken as the scalar component of a chiral superfield, and its potential is derived from the superpotential $I = (\Delta^2/M)(\phi - M)^2$. This is the simplest choice that satisfies the conditions that supersymmetry remains unbroken in the minimum of the potential and that the present cosmological constant is zero. The observed cosmic microwave background anisotropy fixes the inflationary scale around $\lambda \equiv \Delta/M \simeq 10^{-4}$. For the sake of simplicity, we will consider the case in which the gravitino mass is much smaller than the effective mass of the inflaton in this model, $m_{3/2} \ll m_\phi \simeq 10^{-8}M$, and, since the production will take place during a few inflaton oscillations, we will neglect the mass term in the equations. The scalar field potential is given by [11]

$$V(\phi) = e^{|\phi|^2/M^2}\left(\left|\,\frac{\partial I}{\partial\phi} + \frac{\phi^* I}{M^2}\right|^2 - \frac{3|I|^2}{M^2}\right). \tag{24}$$

For the above superpotential, the imaginary direction is known to be stable and therefore we will take for simplicity a real inflaton field. Along the real direction the potential can be written as [13]

$$V(\phi) = \lambda^4 e^{\phi^2}\{[2(\phi - 1) + \phi(\phi - 1)^2]^2 - 3(\phi - 1)^4\}, \tag{25}$$

where we are working in units $M = 1$. We will assume, as indicated in [13], that the potential contributions of dilaton and moduli fields are fixed during and after inflation. This potential has a minimum at $\phi = 1$. The coupling of the inflaton field to gravitinos is given by the following mass term in the supergravity Lagrangian [11]:

$$\mathcal{L} = -\frac{1}{4}\, e^{G/2}\bar\psi_\mu[\gamma^\mu, \gamma^\nu]\psi_\nu\,, \tag{26}$$

$$e^{G/2} = \lambda^2 e^{\phi^2/2}(\phi - 1)^2, \tag{27}$$

where we have chosen the minimal form for the Kähler potential $G(\Phi, \Phi^\dagger) = \Phi^\dagger\Phi + \log|I|^2$. The rest of the interaction terms in the supergravity Lagrangian are not relevant for our purposes. The inflaton and Friedmann equations can be written in conformal time as

$$\ddot\phi + 2\,\frac{\dot b}{b}\,\dot\phi + \frac{b^2}{\lambda^4}\,V_{,\phi} = 0\,, \tag{28}$$

$$\frac{\dot b^2}{b^2} = \frac{1}{3}\left(\frac{1}{2}\,\dot\phi^2 + \frac{b^2}{\lambda^4}V\right), \tag{29}$$

where the derivatives are with respect to the new time coordinate $\tilde\eta = a_0\lambda^2\eta$ and the new scale factor is defined as $b(\tilde\eta) = a(\tilde\eta)/a_0$ with $a_0 = a(0)$. The solution of this equation shows that after the inflationary phase the scalar field starts oscillating around the minimum of the potential with damped amplitude. Substituting in (22) for this particular case we obtain

$$\left(\frac{d^2}{d\tilde\eta^2} + \kappa^2 + \frac{i}{\lambda^2}\,\frac{d}{d\tilde\eta}\,(be^{G/2}) + \frac{b^2}{\lambda^4}\,e^G\right) f_{\kappa l}(\tilde\eta) = 0\,, \tag{30}$$

with $\kappa = p/(a_0\lambda^2)$. From this expression we see that when the scalar interaction is switched off there is no particle production, even in the expanding background. Following [2,14] we can calculate the occupation number:

$$N_{\kappa l}(\tilde T) = \frac{1}{4\kappa}\left(2\kappa + i[\dot f_{\kappa l}^*(\tilde T) f_{\kappa l}(\tilde T) - f_{\kappa l}^*(\tilde T)\dot f_{\kappa l}(\tilde T)] - \frac{2}{\lambda^2}\, be^{G(\tilde T)/2}[|f_{\kappa l}(\tilde T)|^2]\right). \tag{31}$$

In order for the particle number to be well defined, we must evaluate it when the interaction is vanishingly small, that is, for large values of $\tilde{T}$. Here $f_{\kappa l}$ is a solution of Eq. (30) with initial conditions $f_{\kappa l}(0) = 1$ and $\dot{f}_{\kappa l}(0) = -i\kappa$ which corresponds to a plane wave for $\tilde{\eta} \leq 0$. In order to define the initial vacuum at $\tilde{\eta} = 0$, we have taken the inflaton to be at the minimum of the potential at that moment $[\phi(0) = 1]$, which implies $e^{G(\phi=1)/2} = 0$ and $b(0) = 1$. We have chosen $\dot{\phi}(0) = 1.8$ in our numerical computations which corresponds to an initial amplitude of the inflaton oscillations of about $0.06M_p$ and a maximum value of the coupling $e^{G/2}$ of $10^{-10}M_p$.

The results for the spectra in the expanding background are shown in Fig. 1. Note that we have not considered the backreaction effects of the produced particles. In the flat space calculation, we find that broad resonance bands may appear, similar to those in [2,15]. When expansion is taken into account (Fig. 1), the production is reduced by 3–4 orders of magnitude, however the number of particles produced is not negligible. From Fig. 1, we can estimate a lower bound to the total number density of gravitinos of both helicities as

$$n(\eta) = \frac{1}{\pi^2 a^3(\eta)} \int_0^\infty N_{pl} p^2 \, dp = \frac{a_0^3 \lambda^6}{\pi^2 a^3} \int_0^\infty N_{\kappa l} \kappa^2 \, d\kappa \, . \tag{32}$$

Comparing with the number density of a thermal distribution of helicity $\pm 1/2$ gravitinos as estimated in [16] (the helicity $\pm 3/2$ could be even less abundant) for a typical value of the scale factor at the end of inflation [17] $a_0 \simeq 10^{-26}$, the vacuum fluctuation production is suppressed by a factor 10^{10}. The corresponding cosmological consequences have been studied in [3,18]. Comparing with the entropy density today we obtain $n/s \geq 10^{-12}$. This result is 4–5 orders of magnitude larger than the perturbative production during reheating from direct inflaton decay [13] and it could pose compatibility problems with the nucleosynthesis bounds [4,6] for some values of the gravitino mass.

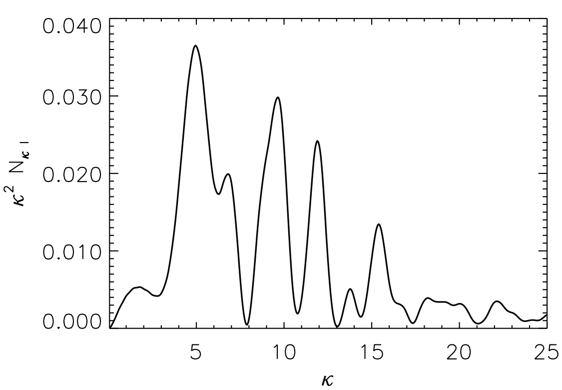

FIG. 1. Number density of helicity $l = \pm 3/2$ gravitinos ($\kappa^2 N_{\kappa l}$) against κ.

We have considered the production of helicity $\pm 3/2$ gravitinos (which are the relevant states for the current nucleosynthesis bounds) in a particular inflationary model. The expression (32) shows that the results are very sensitive to the model parameters, but they can be used to discriminate between different supergravity inflationary models. The completion of the picture would require to study other models and also include the production of helicity $\pm 1/2$ modes; however, the Bogolyubov technique appears very involved for this purpose. (After the appearance of this work, the helicity $\pm 1/2$ case was considered in Kallosh *et al.*, hep-th/9907124.)

A. L. M. acknowledges support from the SEUID-Royal Society and (CICYT-AEN97-1693) (Spain). A. M. is supported by INLAKS and ORS. We thank Andrew Liddle for valuable discussions.

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