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Gravitino-induced baryogenesis: A problem made a virtue

James Cline and Stuart Raby*
Department of Physics, The Ohio State University, Columbus, Ohio 43210
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We investigate cosmic baryon generation in the minimal low-energy supergravity model with the addition of dimension-four baryon-violating interactions. It is shown that gravitinos, if heavier than squarks, can give rise to a baryon asymmetry as they decay. Assuming that gravitinos are not diluted by inflation in the post-Planckian epoch, we find a realistic baryon-to-photon ratio for various ranges of gaugino and squark masses between 100 GeV and 10 TeV. We thus solve two problems simultaneously—the question of whence baryons originate and how to avoid washing out the baryon-photon ratio due to entropy production from decaying gravitinos.

I. INTRODUCTION

In big-bang cosmology a plausible explanation for the observed matter-antimatter asymmetry requires reactions that violate both CP and baryon number (B) and occur out of equilibrium.[1] The first two conditions depend solely on microphysics, while the latter needs a delicate conspiracy of microphysics and macrophysics, and is typically the most difficult to attain. Initial attempts[2] focused on the decays of heavy vector mesons with mass $M_v \geq 10^{-3} M_{\rm Pl}$. This mechanism had the advantage of naturally incorporating out-of-equilibrium decays, where the decay rate for V mesons is much less than the expansion rate at temperatures $T \approx M_v$: $\Gamma_v \approx \alpha M_v \ll H(M_v) \approx M_v^2/M_{\rm Pl}$. Here α, a generic fine-structure constant, is constrained by the out-of-equilibrium condition to be $\alpha \ll M_v/M_{\rm Pl}$, which might be a reasonable value.

Although these attempts were partially successful, they are somewhat disappointing for several reasons. Since the necessary CP violation occurred only in B-violating processes, and it now appears unlikely that we will ever see protons decay, these symmetry-breaking interactions are not directly observable; their only evidence is the asymmetry which they were invented to explain. In the context of inflationary universe scenarios there is an additional problem: after inflation it is difficult to reheat the universe to the high temperatures needed to rethermalize the heavy mesons and yet maintain acceptably small density fluctuations.[3] Finally, it has been realized that the anomalous B-violating interactions in the electroweak theory can wash out any B asymmetry generated at temperatures above M_W/α, unless $B-L\neq 0$.[4,5] This may be a further strong constraint on any high temperature baryogenesis model, unless the anomaly actually *produces* the B asymmetry, as suggested by some authors.[4,6]

Several theories of low temperature baryogenesis exist in the literature.[6–8] The major benefit of these models is their testability, due to possible signs of the necessary B and CP violation in high-energy accelerators and particle electric dipole moments (EDM's). The major obstacle in constructing such models, aside from experimental limits on CP and B violation, is getting the interactions to occur out of thermal equilibrium. At low temperatures the expansion rate is generally much slower than particle reactions. Departure from equilibrium can nevertheless be achieved by the late decay of some heavy particle or by a first-order phase transition.

In the minimal low-energy supergravity (MLES) model, gravitinos decouple at temperatures T somewhat below $M_{\rm Pl}$ and decay very late, at $T \sim 1$ MeV for gravitino masses of order 10 TeV. At this time they dominate the energy density of the universe; thus (1) if they decay after nucleosynthesis, the energetic decay products would dissociate the nascent light elements and (2) the entropy generated in the gravitino decays (typically ΔS of order $10^4 S$) would severely wash out the baryon asymmetry generated prior to their decay. The latter led Weinberg to conclude that $m_{\tilde G} > 10$ TeV.[9] It was suggested that inflation could solve both these problems by diluting the gravitino density.[10] To avoid regenerating the gravitinos after the inflationary epoch, thus preserving the products of nucleosynthesis, it was argued that the reheat temperature must be less than 10^8 GeV m_{100}^{-1} with $m_{100} \equiv m_{\tilde G}/100$ GeV for a gravitino mass in the range 1 GeV $< m_{\tilde G} < 10^4$ GeV.[11] Such a condition makes high-temperature baryogenesis in supersymmetric inflationary scenarios unlikely.

The reader should note, however, that these same gravitinos which are apparently wreaking havoc are out of thermal equilibrium when they decay. Their energy has continued to redshift with the expansion of the universe and their number in a comoving volume remains constant. Thus for temperatures $T > m_{\tilde G}$ their distribution function remains thermal, but not so for $T < m_{\tilde G}$. Hence gravitinos naturally provide the out-of-equilibrium decays needed for baryogenesis. Furthermore, CP violation is generically present in the MLES model, with stringent limits coming from the neutron electric dipole moment. We need only add baryon-violating reactions. Dimopoulos and co-workers have considered the experimental consequences of renormalizable B-violating, lepton-number-conserving interactions that are allowed in any low-energy supergravity model.[12] These interactions do not contribute to proton decay but are constrained by $\Delta B=2$ processes such as n-$\bar{n}$ oscillations or heavy nuclei decays. (Reference 8 also discussed generating a B asymmetry with these interactions, using inflatons, however, as the late-decaying particles.)

We thus have the means to solve two problems simultaneously: no longer is it necessary to find clever ways of disposing of the gravitino, for it now becomes the source of baryonic matter in the cosmos. Here, in contrast with previous supersymmetric cosmological models, we assume that the primordial gravitino density is undiluted by inflation, or sufficiently regenerated afterwards, so that gravitinos dominate the energy density at their decay time. At temperatures of order 1 MeV gravitinos decay and produce many energetic particles. Baryogenesis may then be envisaged as a two-step process. In the first step, the CP-violating decays of gauginos (either those produced by gravitino decays, or the gravitino itself) produce an asymmetry in the number density of squarks and antisquarks. In step two, the subsequent B-violating decays of the squarks convert this asymmetry into a genuine baryon asymmetry. We show that gravitino decays can naturally generate a baryon-to-photon ratio of order 10^{-10} with a gravitino mass of order 10 TeV and squark and other gaugino masses in the range 0.1 to 10 TeV. The model has many testable consequences. Perhaps one of the most sensitive tests will be a measurement of the neutron EDM of order $10^{-26}\,e$ cm, just one order of magnitude below the current experimental limit. Moreover, if squarks are light enough to be seen at the Fermilab Tevatron, Superconducting Super Collider (SSC), or CERN LHC, then their baryon-number-violating decays should be dramatic. For example, a single squark may be produced by a B-violating quark-quark fusion process in pp or $p\bar{p}$ collisions. The subsequent B-violating decay of the squark will be visible as a bump in the two-jet invariant-mass spectrum. Such events should also produce anomalously large numbers of strange particles.[12]

Our paper is organized as follows. We first recount the cosmology of decaying gravitinos, deriving the baryon-to-photon ratio η and reheat temperature T_R; the former depends on ϵ, the net baryon number produced per gravitino decay. We then calculate ϵ from the interference of CP-violating diagrams for gravitinos, and possibly lighter gauginos (which are themselves decay products of the gravitino), decaying into quark-squark pairs. Because of the CP violation, there is an asymmetry between the number of squarks and antisquarks produced. In both cases this is translated into a baryon asymmetry by the fast B-violating decays of t, b, and s squarks. Previous results on supersymmetric (SUSY) contributions to the neutron EDM are used to constrain the CP-breaking phases. We give examples of SUSY masses that lead to the observed baryon asymmetry, all of which include some particles in the 100-GeV region. We conclude with some remarks on the possibility of a nonstandard nucleosynthesis, where baryons are produced below the temperature at which weak interactions establish chemical equilibrium between neutrons and protons.

II. COSMOLOGY WITH DECAYING GRAVITINOS

In the following we will use the sudden decay approximation and assume that the universe is dominated by nonrelativistic gravitinos at the decay time τ, given by

$$\tau^{-1}=\Gamma_{\tilde{G}}=\kappa\frac{m_{\tilde{G}}^3}{M^2}\ , \tag{1}$$

where $M^2=(8\pi G)^{-1}$ and κ is a number of order 1, to be computed below. Equating τ^{-1} with the expansion rate gives the energy density, hence the number density, of gravitinos at the decay time:

$$n_{\tilde{G}}=\frac{\rho_{\tilde{G}}}{m_{\tilde{G}}}=3\kappa^2\frac{m_{\tilde{G}}^5}{M^2}\ . \tag{2}$$

Converting $\rho_{\tilde{G}}$ into g_R species of radiation gives the reheat temperature

$$T_R=\left[\frac{3\kappa^2}{g_R a}\right]^{1/4}\left[\frac{m_{\tilde{G}}^3}{M}\right]^{1/2}\ ,\quad a\equiv\pi^2/30\ . \tag{3}$$

(We assume that the light particles thermalize instantaneously, a good approximation for electrons, positrons, and photons due to their electromagnetic interactions. Whether neutrinos thermalize at all, however, depends sensitively on T_R.) We are interested in reheat temperatures $\gtrsim 1$ MeV so that neutrinos are in thermal equilibrium and the neutron-proton ratio will evolve in the manner of standard nucleosynthesis; then $g_R\simeq 10.75$. The number of photons right after reheating is $n_\gamma=2\zeta(3)T_R^3/\pi^2$ $[\zeta(3)=1.202]$ and increases by a factor of 11/4 during the era of e^+e^- annihilation. If each gravitino on average produces ϵ baryons, the present ratio of baryons to photons will be

$$\eta\simeq\frac{4}{11}\epsilon\frac{n_{\tilde{G}}(\tau)}{n_\gamma(\tau)}=\epsilon g_R\frac{2a\pi^2}{11\zeta(3)}\left[\frac{T_R}{m_{\tilde{G}}}\right]\ . \tag{4}$$

The value of η inferred from nucleosynthesis calculations[13] (about ten times greater than the lower bound from direct observations[14]) is

$$2.6\lesssim\eta_{10}\equiv 10^{10}\eta\lesssim 4.3\ . \tag{5}$$

In the above derivation no mention was made of the "decay temperature," which naively would be reached before reheating. This is because the universe never really heats up due to particle decays;[15] it merely cools more slowly. Therefore T_R is the approximate temperature at the decay time, and the universe does not undergo nucleosynthesis twice, as is sometimes supposed. However Ref. 15 shows that the estimate of the entropy produced in the instantaneous decay approximation is within 10% of the exact result, so we still expect Eq. (4) for η to be quantitatively accurate.

Let us examine under what conditions our assumption holds, that gravitinos dominate the energy density at the decay time. If $\rho_{\tilde{G}}(\tau)$ is *equal* to the energy density of relic radiation $\rho_{\text{relic}}(\tau)$, Eq. (4) overestimates the baryon-photon ratio only by a factor of $2^{3/4}$, so it is not a bad approximation to continue to ignore $\rho_{\rm relic}$ even in this limiting case. Note that in this approximation the temperature T_D of the relic radiation at $t=\tau$ will be identical to T_R, Eq. (3). Denote the ratio of gravitinos to relic photons at time τ by $f\equiv(n_{\tilde G}/n_{\gamma,\rm relic})_{t=\tau}$, or, in terms of energy densities,

$$f=\frac{g_R a\pi^2}{2\zeta(3)}\left[\frac{T_D}{m_{\tilde G}}\right]\left[\frac{\rho_{\tilde G}}{\rho_{\rm relic}}\right]_{t=\tau}, \tag{6}$$

and at first suppose there was no inflation. Were it not for the annihilation of $g_{\rm Pl}-g_R$ heavy species since the Planck era, f would simply be unity, but because of the latter, $f=3g_R/2g_{\rm Pl}$. Using (6), the gravitino-domination condition becomes $g_{\rm Pl}<m_{\tilde G}/T_R$, which is true in any reasonable model (in ours, $m_{\tilde G}/T_R\sim10^7$). However, if inflation has diluted $n_{\tilde G}$ since the Planck time, the story is different: the universe must reheat at the end of inflation to a temperature $T_{\rm max}$ high enough so that gravitinos are adequately regenerated. Reference 11 calculated the regenerated density at $T=T_{\rm max}$ to be $n_{\tilde G}=3.35\times10^{-21}T^4_{\rm max}[1-0.018\ln(T_{\rm max}/10^9\ {\rm GeV})]$ ${\rm GeV}^{-1}$, assuming $T_{\rm max}\ll M_{\rm Pl}$. Thus instead of $3g_R/2g_{\rm Pl}$, one finds

$$f=3.35\times10^{-21}\left[\frac{T_{\rm max}}{1\ {\rm GeV}}\right]\frac{3\pi^2}{4\zeta(3)}\times\left[1-0.018\ln\left[\frac{T_{\rm max}}{10^9\ {\rm GeV}}\right]\right](g_R/g_{\rm Pl})\,. \tag{7}$$

Again using (6), the gravitino-domination condition implies a lower limit for $T_{\rm max}$:

$$T_{\rm max}>2\times10^{15}\ {\rm GeV}\left[\frac{T_R}{1\ {\rm MeV}}\right]\left[\frac{m_{\tilde G}}{10\ {\rm TeV}}\right]^{-1}\left[\frac{g_{\rm Pl}}{200}\right], \tag{8}$$

where we have anticipated the values of the various parameters that will be of interest in the remainder of the paper. (Ironically, Ref. 11 used their estimate of $n_{\tilde G}$ to get an upper limit on $T_{\rm max}$.) Notice that f and, hence $n_{\tilde G}$, is much smaller in the inflationary case than otherwise; however this will not affect the baryon-photon ratio as long as gravitinos dominate at $t=\tau$. This is simply because the number of photons produced by the decays is proportional to the number of gravitinos decaying, and these photons outnumber the relic photons if $\rho_{\tilde G}>\rho_{\rm relic}$.

III. GENERATION OF THE BARYON ASYMMETRY

Our particle-physics model is the simplest possible supergravity theory that allows for baryon violation and a large splitting between $m_{\tilde G}$ and the squark masses ($\tilde m$). We assume the minimal particle content and impose lepton-number conservation. The lightest supersymmetric particle (LSP) is unstable because of the B-violating interactions. These come from a term in the superpotential,[12]

$$g_{ijk}U_i^cD_j^cD_k^c\ , \tag{9}$$

involving SU(2)-singlet superfields U^c, D^c, containing up- and down-type quarks, respectively, with generation indices i,j,k. To prevent neutron oscillations and heavy nuclei decay, g_{112} must be small, $\lesssim10^{-6}(\tilde m/300\ {\rm GeV})^{5/2}$,[8] whereas some of the g_{ijk} must be larger to produce a sufficient baryon asymmetry. For simplicity (and to keep our estimates conservative) we will take only g_{332} to be nonvanishing. This coupling induces radiative corrections to g_{112}-like operators via the diagram in Fig. 1, which we estimate to be

$$g_{332}(\alpha_W/4\pi)V_{ub}V_{td}m_bm_t(M_W^2-m_t^2)^{-1}\ln(M_W^2/m_t^2)\sim10^{-8}g_{332}$$

for a top-quark mass of 100 GeV. Using the above limit on g_{112}, we have

$$\alpha_{332}\equiv\frac{g_{332}^2}{4\pi}\lesssim3.3\left[\frac{\tilde m}{100\ {\rm GeV}}\right]^5; \tag{10}$$

i.e., α_{332} is practically unconstrained. The same calculation with charmed instead of top quarks gives an induced coupling that is only seven times greater. Thus we could also take α_{223} to be ~0.5 and still have a naturally small value for g_{112}.

In addition to B violation, CP violation is required to produce the baryon asymmetry.[1] This arises naturally in supergravity from several independent phases, including that of the parameter A in the supersymmetry-breaking part of the low-energy effective potential. (The MLES model has five CP-breaking phases, in addition to the less effectual Kobayashi-Maskawa matrix phases. In the approximation we use of quarks being much lighter than supersymmetric particles, only four of these contribute to processes that generate a baryon asymmetry: ${\rm Im}\,Am_{\tilde x}$, where $\tilde x=\tilde G,\tilde g,\tilde Z$, and $\tilde\gamma$.) In the case of a U(N)-invariant Kähler potential, the latter takes the special form[16]

$$\mathcal{L}_{\rm sym\ breaking}=\tilde m^2\sum_i|z_i|^2+AW^{(3)}+(A-m_{\tilde G})W^{(2)}+(A-2m_{\tilde G})W^{(1)}\ , \tag{11}$$

where z_i are the scalar fields, $W^{(n)}$ is the dimension n, purely bosonic part of the superpotential, and $|A|$ is some low-energy, SUSY-breaking mass scale. Equations

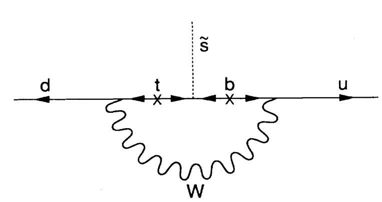

FIG. 1. Radiative contribution to an operator that could cause heavy nuclei to decay.

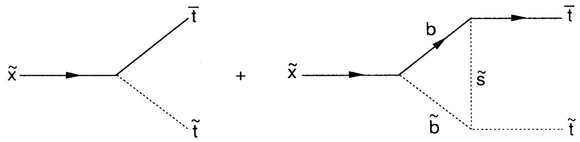

FIG. 2. Example of diagrams whose interference produces an asymmetry between the density of squarks and antisquarks in $\tilde{x}$-gaugino decays.

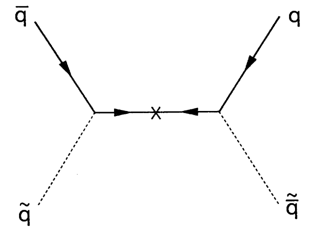

FIG. 3. Scattering process that would erase the CP asymmetry produced in Fig. 2, but is highly suppressed by low particle densities.

(9) and (11) thus give rise to the B- and CP-violating interactions

$$\mathcal{L}_{B,CP\ \text{violating}} = A g_{332}\tilde{t}^c\tilde{b}^c\tilde{s}^c - g_{332}(\bar{T}P_+B^c\tilde{s}^c + \bar{B}P_+S^c\tilde{t}^c + \bar{S}P_+T^c\tilde{b}^c) + \text{H.c.}\ , \tag{12}$$

where upper (lower) case letters are quark (squark) fields, $P_+ = (1+\gamma_5)/2$, and antisymmetrization with respect to color is implied.

Depending on the sign of $\mathrm{Im}(Am_{\text{gaugino}})$, gauge fermions, including the gravitino, decay preferentially into squarks or antisquarks through the interference of diagrams such as in Fig. 2. If all gauginos are heavier than squarks, the only tree diagrams through which the latter can decay are the $\Delta B = 1$ interactions of Eq. (12). Therefore the asymmetry between $\tilde{s}^c, \tilde{b}^c, \tilde{t}^c$ antisquarks and squarks will be at least partially converted into a baryon asymmetry, at a rate of $\frac{1}{2}\alpha_{332}\tilde{m}$. This is many orders of magnitude faster than the rate at which the CP asymmetry is washed out by the rescatterings shown in Fig. 3; these events are suppressed by the low particle density [Eq. (2)] at the decay time. In view of the above, we can write an expression for ϵ, the net baryon number produced per gravitino decay. Let $\Delta\Gamma_{\tilde{x}}$ denote the rate for $\tilde{x}$ gauginos to generate a CP asymmetry, i.e., $\Delta\Gamma_{\tilde{x}} = \Gamma_{\tilde{x}\to q\tilde{q}} - \Gamma_{\tilde{x}\to\bar{q}\bar{\tilde{q}}}$; let $\Gamma_{\tilde{G}\to\tilde{x}}$ be the partial width for $\tilde{x}$ production, $\Gamma_{\tilde{q}\to\Delta B}$ that for squarks to have $\Delta B = 1$ decays, and $\Gamma_{\tilde{x}}$ the total $\tilde{x}$ decay rate. Then, in the narrow-width approximation,

$$\epsilon = \left[\frac{\Delta\Gamma_{\tilde{G}}}{\Gamma_{\tilde{G}}} + \sum_{\tilde{x}=\tilde{g},\tilde{Z},\tilde{\gamma}} \frac{\Gamma_{\tilde{G}\to\tilde{x}}}{\Gamma_{\tilde{G}}}\frac{\Delta\Gamma_{\tilde{x}}}{\Gamma_{\tilde{x}}}\right]\frac{\Gamma_{\tilde{q}\to\Delta B}}{\Gamma_{\tilde{q}}}\ . \tag{13}$$

Evaluating (13) is greatly simplified if we assume quarks are much lighter than squarks. Although not necessarily true for the top quark, we find that the relative size of the corrections to this approximation are of order $m_t^2|A|/\max(|m_{\tilde{x}}|,\tilde{m})^3 \lesssim 1$. In any case, the same analysis applies to the g_{223} couplings, and there the charmed-quark mass *is* negligible compared to $\tilde{m}$. In the massless quark limit, the squark fields in (12) are mass eigenstates, and the W-ino cannot couple to them since they are SU(2) singlets: the second diagram in Fig. 2 is helicity suppressed. Computing these diagrams as well as the total widths for photino, Z-ino, gluino, and gravitino decay gives the relative CP-violation rates

$$\frac{\Delta\Gamma_{\tilde{x}}}{\Gamma_{\tilde{x}}} = \alpha_{332}\frac{\mathrm{Im}(Am_{\tilde{x}})}{|m_{\tilde{x}}|^2}\begin{cases}\frac{1}{8}f_1(R_{\tilde{q},\tilde{\gamma}}), & \tilde{x}=\tilde{\gamma},\\ \frac{1}{8}(1+\frac{7}{18}\sin^4\theta_W)^{-1}f_1(R_{\tilde{q},\tilde{z}}), & \tilde{x}=\tilde{Z},\\ \frac{1}{4}f_1(R_{\tilde{q},\tilde{g}}), & \tilde{x}=\tilde{g},\\ -\dfrac{3}{32\pi\kappa}f_2(R_{\tilde{q},\tilde{G}}), & \tilde{x}=\tilde{G}\ ,\end{cases} \tag{14}$$

where $R_{\tilde{q},\tilde{x}} = \tilde{m}^2/|m_{\tilde{x}}|^2$,

$$\begin{aligned} f_1(R) &= \theta(1-R)[1-R(1-R)^{-2}\ln(R+R^{-1}-1)]\ ,\\ f_2(R) &= \theta(1-R)\{(1-R)^4+6R(1-R)^2\\ &\quad -[4R(1-R)^2+6R^2]\\ &\quad \times\ln(R+R^{-1}-1)\}\ , \end{aligned} \tag{15}$$

and for simplicity we have neglected mixing of the electroweak-inos (valid if they are much heavier than M_W). The total widths in (14) and the branching ratios for $\tilde{G}\to\tilde{x}$ were calculated to lowest order in the gauge couplings and M_{Pl}. Denoting the dimension of the adjoint representation of the gauge group corresponding to $\tilde{x}$ gauginos by $\mathcal{C}_{\tilde{x}}$, the branching ratios are

$$\Gamma_{\tilde{G}\to\tilde{x}}/\Gamma_{\tilde{G}} = (8\pi\kappa)^{-1}\mathcal{C}_{\tilde{x}}(1-R_{\tilde{x},\tilde{G}})^3(1+R_{\tilde{x},\tilde{G}}/3)\ . \tag{16}$$

In the same approximation, the total width parameter κ [Eq. (1)] is given by

$$8\pi\kappa=\sum_{\tilde{x}} \mathscr{C}_{\tilde{x}}(1-R_{\tilde{x},\tilde{G}})^3(1+R_{\tilde{x},\tilde{G}}/3) + \tfrac{1}{12}\mathscr{N}(1-R_{\tilde{q},\tilde{G}})^4\,, \tag{17}$$

where $\mathscr{N}$ is the number of chiral fields and for simplicity we have taken squarks, sleptons, and Higgs bosons to have the common mass $\tilde{m}$. In minimal SUSY, $\mathscr{N}=36+9+4=49$, and (17) implies $\kappa\leq 0.64$, with equality being reached in the $m_{\tilde{G}}\to\infty$ limit. [Using this value for κ in (3), together with the minimum reheat temperature needed for nucleosynthesis, $T_R\geq 0.4$ MeV,[9] we find that $m_{\tilde{G}}\gtrsim 8$ TeV, close to Weinberg's estimate $m_{\tilde{G}}\gtrsim 10$ TeV.] The last factor in (13), the branching ratio for squarks to violate B, is

$$\Gamma_{\tilde{q}\to\Delta B}/\Gamma_{\tilde{q}}=\left[1+\frac{4\alpha_s}{3\alpha_{332}}(1-R_{\tilde{g},\tilde{q}})^2\theta(1-R_{\tilde{g},\tilde{q}})\right]^{-1}\,, \tag{18}$$

neglecting electromagnetic decays.

IV. EXPERIMENTAL CONSTRAINTS

The CP-violating phases in Eq. (14) are constrained by their contributions to the neutron electric dipole moment d_n. Loops with gauginos and squarks give each quark a chromoelectric dipole moment (CDM) as well as an EDM proportional to $\mathrm{Im}\,Am_{\tilde{x}}$. Running them down from $\Lambda=\max(|m_{\tilde{g}}|,|m_{\tilde{\gamma}}|)$ to the chiral-symmetry-breaking scale $M_{\rm CSB}=1.2$ GeV, they are, respectively,[17,18]

$$\frac{d_q^{\gamma}}{e}=C_{\rm QCD}m_q(M_{\rm CSB})\left[\mathrm{Im}(A_q m_{\tilde{g}})\frac{\alpha_s(\Lambda)Q_q}{3\pi|m_{\tilde{g}}|^4}f_3'(R_{\tilde{m},\tilde{g}})+\mathrm{Im}(A_q m_{\tilde{\gamma}})\frac{\alpha Q_q^3}{4\pi|m_{\tilde{\gamma}}|^4}f_3'(R_{\tilde{m},\tilde{\gamma}})\right]\,, \tag{19a}$$

$$\frac{d_q^{g}}{g_s}=\sqrt{C_{\rm QCD}}\,m_q(M_{\rm CSB})\mathrm{Im}(A_q m_{\tilde{g}})\frac{\alpha_s(\Lambda)}{3\pi|m_{\tilde{g}}|^4}\left[f_3'(R_{\tilde{m},\tilde{g}})+\tfrac{9}{8}f_4'(R_{\tilde{m},\tilde{g}})\right]\,, \tag{19b}$$

where the QCD correction factor is $C_{QCD}\simeq 0.8$ for the mass ranges we are interested in, Q_q is the quark charge and

$$f_3(R)=\frac{-2R\ln R+R^2-1}{2(1-R)^3}\,,$$
$$f_4(R)=\frac{R\ln R+1-R}{(1-R)^2},\quad f'\equiv\frac{df}{dR}\,. \tag{20}$$

[To stem the tide of formulas, we computed only the gluino and photino contributions to (19a); the latter becomes competitive only for $m_{\tilde{g}}/m_{\tilde{\gamma}}$ very large, but we will be pushing $m_{\tilde{g}}$ to rather large values to suppress the neutron EDM in what follows, so it was necessary to verify that the photino contribution did not start to dominate. $\tilde{W},\tilde{Z}$, and Higgsino masses will henceforth be kept large enough to justify ignoring their effects. Equation (19b) is taken from Ref. 18, which also computes the gluino contribution to (19a). We agree with Ref. 18 (but not Ref. 17) on the overall normalization.] A_q in (19) denotes the combination $A\pm\mu(\langle\bar{H}\rangle/\langle H\rangle)^{\pm 1}$, where μ multiplies the $H\bar{H}$ Higgs term in the superpotential, and the sign depends on whether the quark is up or down type. (The phase of μ can be taken to be the remaining CP-violating parameter in MLES.) In order to relate the phases in (19) to those in (14), we will assume there are no accidental cancellations of imaginary parts; explicitly, that $\mathrm{Im}(A_q m_{\tilde{x}})\approx\mathrm{Im}(Am_{\tilde{x}})\approx\mathrm{Im}\,A|m_{\tilde{x}}|$.

The quark dipole moments (19a) and (19b) give rise to a neutron EDM (d_n) through valence-quark contributions

$$\frac{4}{3}d_d^{\gamma}-\frac{1}{3}d_u^{\gamma}+\frac{e}{\sqrt{6}}\sum_{q=u,d}\frac{d_q^g}{g_s}\,, \tag{21}$$

and through heavy quark loops that induce CP-odd multigluon (-photon) operators at the scale $M_{\rm CSB}$, whose respective contributions to d_n are

$$G G\tilde{G}:\ \frac{e}{3\sqrt{6}}\sum_{q=s,c,b,t}\frac{M_{\rm CSB}}{m_q}f(q)^{-54/(33-2N_f)}\frac{d_q^g}{g_s}\,, \tag{22a}$$

$$GGG\tilde{F}+FGG\tilde{G}:\ \frac{e}{12\sqrt{6}}\sum_{q=s,c,b,t}\left[\frac{M_{\rm CSB}}{m_q}\right]^3\times f(q)^{-37.5/(33-2N_f)}\times\left[3Q_q\frac{d_q^g}{g_s}+\frac{d_q^{\gamma}}{e}\right]\,,$$

where $f(q)=\frac{1}{9}(33-2N_f)\ln(m_q/150\text{ MeV})$ and N_f is the number of flavors lighter than m_q.[18] Equations (19)–(22) can be used to relate $\mathrm{Im}\,A$ to d_n. The terms preceded by summation signs are estimated using dimensional analysis; since this does not predict their relative signs, we choose to add them more equitably by computing the root mean square. [This nicety is not crucial for the examples given below since, in most of them, the d^{γ} terms from (21) dominate.] It is doubtful whether strange quarks are light enough to safely integrate out at the scale $M_{\rm CSB}$, so we will compute everything two ways, both keeping and omitting the $q=s$ terms in (22). This gives an indication of the theoretical uncertainties. Also uncertain is the possible direct contribution of strange sea quarks to d_n, which we have therefore ignored.

V. RESULTS AND CONCLUSIONS

Equations (3) and (4) and (13)–(22) contain everything needed to evaluate the baryon asymmetry in our model. The free parameters are the superpartner masses, the B-violating couplings g_{ijk}, and the CP-violating phase of A, which we have traded for the EDM of the neutron. To reduce the dimensionality of this parameter space, we take $\alpha_{332}=0.1$ (small enough for perturbation theory to apply) and $|d_n|=10^{-26}\,e\,\mathrm{cm}$, which if correct will be detected in the next generation of experiments. [The most recent measurements of d_n are $(-1.4\pm0.6)\times10^{-25}\,e\,\mathrm{cm}$ (Ref. 19) and $(-0.3\pm0.5)\times10^{-25}\,e\,\mathrm{cm}$ (Ref. 20), the latter giving a 95%-C.L. limit of $|d_n|<1.2\times10^{-25}\,e\,\mathrm{cm}$.] For different values of these quantities, the following results for the baryon asymmetry η simply scale linearly. As noted previously, in order to have $T_R>0.4$ MeV, $m_{\tilde G}$ is constrained to be greater than $\sim 8-10$ TeV.

We present T_R and η_{10} for some possible values of the gaugino and squark masses in Table I. The results fall roughly into three qualitatively different regimes for the superpartner masses that yield a sufficiently large baryon asymmetry. (a) In the first case, squarks and sleptons are the LSP's. Table I(a) gives examples: the gauge fermions are several TeV and $\tilde m\sim100$ GeV. One finds that η is an increasing function of all the masses (but depends rather weakly on $m_{\tilde z}$). The baryon asymmetry is due primarily to gluino decays, with the photinos making up 25% of the total in the $m_{\tilde\gamma}=1$ TeV case. The heavy-quark contribution (22) to d_n is always less than that of the valence quarks (21) in these examples. There is a noticeable drop in $\eta\propto g_R^{3/4}$ when $m_{\tilde G}$ falls below ~17 TeV, since T_R becomes too small to have neutrinos in thermal equilibrium, and the number of relativistic species goes from 10.75 to 5.5. (b) In the second case the scalars are heavier ($\sim$0.5–5 TeV) than the photinos ($\sim$100 GeV), as exemplified in Table I(b). $m_{\tilde g}$ must be $\gtrsim3$ TeV or d_n would start to exceed $10^{-26}e$ cm. Gluino decays once again dominate in producing the asymmetry in all but the last case, $\tilde m=4$ TeV, where direct decays of gravitinos into squark-quark pairs is the only possibility. (c) The gravitino can be made light if we relax the requirement that T_R exceed 0.4 MeV. In this case the neutron-to-proton ratio is not kept in equilibrium by weak interactions, so nucleosynthesis would not proceed in the usual manner, see the discussion below. (Actually 0.8 MeV may be a more accurate estimate for the $n\leftrightarrow p$ freeze-out temperature.[21])

In conclusion, we have shown that late decaying gravitinos can naturally generate the observed cosmological baryon asymmetry. We have given examples in which the neutron EDM and either the scalar or the fermionic superpartners are just on the verge of detectability. However we have been forced to keep the gravitino mass $\gtrsim10$ TeV so that gravitino decays will reheat the universe to a temperature above 0.4 MeV. This assures that (1) neutrinos, hence neutrons and protons, will be in chemical equilibrium; (2) the decays occur early enough in the history of the universe so very energetic decay products will not cause photofission of the light elements;[22] (3) in general the density of decay products is sufficient for them to annihilate before nucleosynthesis begins. (For example, we find numerically that the ratio of antiprotons to photons is $\sim10^{-11}$ when $T=0.03$ MeV for 2-TeV gravitinos, assuming each $\tilde G\to1000p\bar p+\cdots$, whereas it is completely negligible for $m_{\tilde G}=20$ TeV. In the former case, the eventual $p\bar p$ annihilations would lead to some photodissociation of nuclei.) It would be desirable to relax this constraint on T_R since a lower $m_{\tilde G}$ would be more natural when considering a lower scale of supersymmetry breaking. As shown in Table I(c), there is no difficulty in making enough baryons with smaller values of $m_{\tilde G}$; however, for $m_{\tilde G}\lesssim2$ TeV, T_R is too low for nuclei to fuse. Above this threshold, the challenge is to explain why the initial conditions for the neutron-proton ratio are just right to produce the observed abundance of helium, which happens naturally in the orthodox version of nucleosynthesis. It is possible that the energetic neutrinos produced in $\tilde G$ decays equilibrate $n\leftrightarrow p$ without heating the photons, or that a suitable ratio occurs fortuitously for some range of superpartner masses. We are numerically exploring the feasibility of nucleosynthesis in the presence of decaying gravitinos, at these lower temperatures.[23]

TABLE I. Reheat temperature (T_R) in MeV and baryon-to-photon ratio ($\eta_{10}=10^{10}\eta$) for different values of the gravitino, gluino, Z-ino, photino, and squark/slepton masses, given in TeV. The value of η_{10} in parentheses is obtained by omitting the dubious (Ref. 18) s-quark contribution to the neutron EDM in Eq. (22); this gives some estimate of the theoretical uncertainty. Values are for $d_n=10^{-26}e$ cm and $\alpha_{332}=0.1$.

	$m_{\tilde G}$	$m_{\tilde g}$	$m_{\tilde z}$	$m_{\tilde\gamma}$	$\tilde m$	T_R	η_{10} (no s)
(a)	20	5	1	2	0.3	1.33	4.1 (4.5)
	20	5	1	2	0.1	1.33	2.4 (2.8)
	20	6	1	2	0.1	1.31	3.0 (3.5)
	20	6	1	1	0.1	1.31	2.0 (2.3)
	15	6	1	2	0.3	0.96	2.2 (2.4)
(b)	20	3.1	1	0.1	0.5	1.37	2.5 (2.7)
	20	4	0.1	0.1	0.5	1.35	3.1 (3.2)
	20	4	0.1	0.1	0.4	1.35	2.3 (2.4)
	20	4	0.1	0.1	4	1.33	1.3 (2.9)
(c)	8	10	8	8	0.5	0.27	2.6 (2.8)
	3	10	9	9	0.3	0.06	3.0 (3.4)

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*On leave of absence from Theoretical Division, Los Alamos National Laboratory, Los Alamos, NM 87545.

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