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Authors

Buchmüller, W
Hamaguchi, K
Ratz, M

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Gauge couplings at high temperature and the relic gravitino abundance

W. Buchmüller, K. Hamaguchi, M. Ratz

Deutsches Elektronen-Synchrotron DESY, Hamburg, Germany

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Abstract

In higher-dimensional supersymmetric theories gauge couplings of the effective four-dimensional theory are determined by expectation values of scalar fields. We find that at temperatures above a critical temperature T_* , which depends on the supersymmetry breaking mass scales, gauge couplings decrease like $T^{-\alpha}$, $\alpha > 1$. This has important cosmological consequences. In particular it leads to a relic gravitino density which becomes independent of the reheating temperature for $T_R > T_*$. For small gravitino masses, $m_{3/2} \ll m_{\tilde{g}}$, the mass density of stable gravitinos is essentially determined by the gluino mass. The observed value of cold dark matter, $\Omega_{\text{CDM}} h^2 \sim 0.1$, is obtained for gluino masses $m_{\tilde{g}} = \mathcal{O}(1 \text{ TeV})$.

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In higher-dimensional supersymmetric theories [1], where the standard model emerges as low-energy effective theory, gauge and Yukawa couplings are determined by expectation values of gauge singlet ‘moduli’ fields. In a cosmological context, this implies that generically all couplings depend on the parameters of the cosmological evolution, such as the Hubble parameter, temperature, or the cosmological constant.

In the following we study the dependence of gauge couplings on temperature. As we shall see, this has important consequences for the production of gravitinos in the early universe. ‘Vacuum alignment’ at high temperatures causes a power-like decrease of gauge couplings. This then leads to a relic gravitino density which becomes independent of the reheating temperature T_R above a critical temperature T_* .

As a specific example, consider gaugino mediation [2,3] which is an attractive mechanism to generate a realistic mass spectrum of gauginos, higgsinos and scalar quarks and leptons in the supersymmetric standard model. The source of supersymmetry breaking is the vacuum expectation value of a gauge singlet chiral superfield S ,

$$\langle S \rangle = S_0 + \theta \theta F_S, \quad (1)$$

which is localized on a four-dimensional (4d) brane embedded in D -dimensional spacetime. The coupling to bulk gauge fields, expressed in terms of 4d $N = 1$ superfields, is given by

$$I_D = \int d^4x d^{D-4}y d^2\theta \left\{ \frac{1}{4g_D^2} W^a W^a + \delta^{(D-4)}(y - y_S) \frac{1}{4M} S W^a W^a + \dots \right\} + \text{h.c.}, \quad (2)$$

E-mail address: buchmuwi@mail.desy.de (W. Buchmüller).

where W^a is the supersymmetric field strength. M is a mass scale in the range between the compactification scale and the D -dimensional Planck mass,

$$\frac{1}{V^{1/(D-4)}} < M < M_D < M_P. \quad (3)$$

Here $V = \int d^{D-4}y$ is the volume of the compact dimensions, $M_D = (VM_D^{D-4})^{-1/2}M_P$ and $M_P = (8\pi G_N)^{-1/2} = 2.4 \times 10^{18}$ GeV is the 4d Planck mass. For instance, with $1/V^{1/(D-4)} \simeq M_{\text{GUT}} = 2 \times 10^{16}$ GeV one obtains $M_D = 2 \times 10^{17}$ GeV in the case $D = 6$.

Inserting the expectation value (1) in the action (2) one obtains for the 4d gauge coupling and for the gaugino mass,

$$\frac{V}{g_D^2} + \frac{\phi_0}{M} = \frac{1}{g_0^2}, \quad (4)$$

$$m_{\tilde{g}} = \frac{g_0^2}{2} \frac{F_S}{M}, \quad (5)$$

where $\phi_0 = \text{Re } S_0$. For the SU(3) gauge coupling of the standard model one has $g_0^2(\mu) \geq g_0^2(M_{\text{GUT}}) \simeq 1/2$. The gravitino mass is given by

$$m_{3/2} = \eta \frac{F_S}{M_P}, \quad (6)$$

where $\eta \geq 1/\sqrt{3}$. The smallest gravitino mass is obtained if F_S is the only source of supersymmetry breaking, which is the case in gaugino mediation. The gravitino mass is then always smaller than the gaugino mass $m_{\tilde{g}}$,

$$m_{\tilde{g}} \geq \frac{g_0^2}{2} \frac{F_S}{M_D} = \frac{\sqrt{3}}{2} g_0^2 (VM_D^{D-4})^{1/2} m_{3/2} > m_{3/2}, \quad (7)$$

since the volume enhancement factor $\rho = (VM_D^{D-4})^{1/2}$ is larger than $2/(\sqrt{3}g_0^2) \leq 4/\sqrt{3}$. For instance, in $D = 6$ one has $\rho \sim 10$.

The 4d effective action for the zero modes contains a coupling of the scalar field ϕ to the supersymmetric gauge kinetic term,

$$I_4 = \int d^4x \left\{ -\frac{1}{4} F_{\mu\nu}^a F^{a\mu\nu} - i\lambda^a \sigma^\mu (D_\mu \bar{\lambda})^a - \frac{1}{2} m_{\tilde{g}} (\lambda^a \lambda^a + \bar{\lambda}^a \bar{\lambda}^a) \right.$$

$$\left. + g_0^2 \frac{\phi}{M} \left(-\frac{1}{4} F_{\mu\nu}^a F^{a\mu\nu} - i\lambda^a \sigma^\mu (D_\mu \bar{\lambda})^a \right) + \dots \right\}; \quad (8)$$

here F^a is the field strength of the vector potential A^a , and λ^a denotes the gaugino. At finite temperature the gauge kinetic term acquires an expectation value which leads to a force on the scalar field ϕ . This expectation value can be easily calculated by making use of the anomalous divergence of the supercurrent [4],

$$\bar{D}^{\dot{\alpha}} J_{\alpha\dot{\alpha}} = \frac{1}{3} \frac{\beta(g_0)}{g_0} D_\alpha W^a W^a, \quad (9)$$

which contains the trace anomaly of the energy-momentum tensor,

$$T^\mu{}_\mu = -2 \frac{\beta(g_0)}{g_0} \left(-\frac{1}{4} F_a^{\mu\nu} F^{a\mu\nu} - i\lambda^a \sigma^\mu (D_\mu \bar{\lambda})^a \right), \quad (10)$$

where $\beta(g_0)$ is the usual β -function of the gauge coupling.

The thermal average of the energy-momentum tensor is determined by energy density and pressure,

$$\langle T^\mu{}_\mu \rangle_T = \varepsilon - 3P, \quad (11)$$

which are related by

$$\varepsilon = -P + Ts = -P + T \frac{\partial P}{\partial T}. \quad (12)$$

The pressure has been calculated in perturbation theory for a gauge theory with fermions in the fundamental representation [5]. Correcting for the colour charge of the gauginos one obtains for a pure supersymmetric gauge theory,

$$P = (a_0 - a_2 g_0^2(T) + \dots) T^4, \quad (13)$$

with

$$a_0 = \frac{\pi^2}{24} n_A, \quad a_2 = \frac{1}{64} T_A n_A. \quad (14)$$

Here T_A is the Dynkin index of the adjoint representation and $n_A = \dim G$, i.e., the number of gluons. For SU(N) one has $T_A = N$ and $n_A = N^2 - 1$. From Eqs. (10)–(13) one obtains for the thermal expectation value of the gauge kinetic term,

$$\left\langle -\frac{1}{4} F_{\mu\nu}^a F^{a\mu\nu} - i\lambda^a (\sigma^\mu D_\mu \bar{\lambda})^a \right\rangle_T = a_2 g_0^2 T^4. \quad (15)$$

Note that the sign of the expectation value is positive and that there is no dependence on the β -function. Because of the anomaly one no longer has $P = \varepsilon/3$.

The mass of a chiral superfield, whose vacuum expectation value breaks supersymmetry, is generally controlled by the supersymmetry breaking mass scale, i.e., $m_\phi \propto m_{3/2}$. Small fluctuations around the minimum are then described by the lagrangian (cf. (8), (15)),

$$\mathcal{L} = \frac{1}{2}(\partial\phi)^2 - \frac{\xi}{2}m_{3/2}^2\phi^2 + a_2g_0^4T^4\frac{\phi}{M}. \quad (16)$$

Hence, the thermal fluctuations of gauge bosons and gauginos induce a negative linear term in the effective potential for ϕ . In many models the parameter ξ is $\mathcal{O}(1)$.

The negative linear term in the effective potential leads to an increase of the field ϕ . Its equilibrium value at finite temperature is given by

$$\phi_T = \frac{a_2g_0^4}{\xi} \frac{T^4}{m_{3/2}^2M}. \quad (17)$$

Note that the fluctuations of ϕ are not in thermal equilibrium and that ϕ does not acquire a thermal mass. According to (4) the shift in ϕ changes the gauge coupling to $g(\phi_T)$,

$$\frac{1}{g_0^2} + \frac{\phi_T}{M} = \frac{1}{g^2(\phi_T)}. \quad (18)$$

This change of the gauge coupling becomes significant at a temperature T_* where $\phi_T/M \sim 1/g_0^2$, i.e.,

$$T_* = \left(\frac{\xi}{a_2g_0^6}\right)^{1/4} (m_{3/2}M)^{1/2}. \quad (19)$$

Here we have assumed that F_S does not depend on temperature, as in the Polonyi model. Using Eqs. (5) and (6) the mass scale M can be expressed in terms of gaugino and gravitino masses, which yields

$$T_* = \left(\frac{\xi}{a_2g_0^2\eta^2}\right)^{1/4} \left(\frac{m_{3/2}^2M_P}{2m_{\tilde{g}}}\right)^{1/2}. \quad (20)$$

Extrapolating Eqs. (17) and (18) to temperatures larger than T_* leads to a rapid decrease of the gauge coupling as $g^2(\phi_T) \propto 1/T^4$.

However, at large values of ϕ_T/M the effective lagrangian (16) is no longer appropriate. First, the decrease of the gauge coupling reduces the force of the

thermal bath on the field ϕ . This backreaction can be taken into account by using as effective potential the free energy density of the thermal system evaluated with the field-dependent gauge coupling,

$$f = -P = (-a_0 + a_2g^2(T, \phi) + \dots), \quad (21)$$

where $g(T, \phi)$ has to be determined from the equations of motion. Second, for large values of ϕ , higher powers of ϕ/M have to be taken into account. This leads to the effective lagrangian

$$\bar{\mathcal{L}} = \frac{1}{2}(\partial\phi)^2 - \frac{1}{2}m_{3/2}^2h(\phi) - a_2g^2(T, \phi)T^4, \quad (22)$$

where

$$g^2(T, \phi) = \frac{g_0^2(T)}{1 + g_0^2(T)k(\phi)}, \quad (23)$$

$$h(\phi) = \xi\phi^2 \left(1 + \mathcal{O}\left(\frac{\phi}{M}\right)\right), \quad (24)$$

$$k(\phi) = \frac{\phi}{M} \left(1 + \mathcal{O}\left(\frac{\phi}{M}\right)\right).$$

$k(\phi)$ replaces the linear term ϕ/M in Eq. (8). The equilibrium value of ϕ is now determined by the equation

$$\frac{h'(\phi_T)(1 + g_0^2k(\phi_T))^2}{k'(\phi_T)} = 2a_2g_0^4 \frac{T^4}{m_{3/2}^2}. \quad (25)$$

For small values of ϕ one recovers Eq. (17). Neglecting corrections $\mathcal{O}(\phi/M)$ for $h(\phi)$ and $k(\phi)$, keeping only the effect of the back reaction, one obtains at large temperatures $\phi_T \propto T^{4/3}$ and correspondingly for the gauge coupling $g^2(T, \phi_T) \propto T^{-4/3}$. This decrease with temperature is much weaker than the T^{-4} fall-off obtained in the linear approximation. We expect that the true decrease, which is determined by the back reaction together with the behaviour of h and k at large values of ϕ , lies somewhere in between.

The time evolution of the field ϕ is determined by the equation of motion

$$\ddot{\phi} + 3H\dot{\phi} + \frac{1}{2}m_{3/2}^2h'(\phi) - \frac{a_2g_0^4}{(1 + g_0^2k(\phi))^2}k'(\phi)T^4 = 0, \quad (26)$$

where H is the Hubble parameter. For $H > m_{3/2}$ the motion is damped whereas for $H < m_{3/2}$ the field ϕ oscillates. During the period of reheating the Hubble parameter generally depends not only on the thermal bath, but also on the time evolution of other fields, in particular the inflaton. The detailed analysis of the time evolution of ϕ is beyond the scope of this Letter. In the following we shall assume that at the end of reheating thermal equilibrium is achieved and that, to good approximation, ϕ is close to its equilibrium value ϕ_T .

The power-like fall-off of gauge couplings at high temperature, $g^2 \propto T^{-\alpha}$ with $\alpha > 1$, has important cosmological implications. An immediate consequence is that one loses thermal equilibrium at a temperature T_{eq} much below the unification scale M_{GUT} . For instance, for $\alpha = 2$, $m_{\tilde{g}} \simeq 1$ TeV and $m_{3/2} \simeq 100$ GeV, one obtains $\Gamma(T_{\text{eq}}) \simeq H(T_{\text{eq}})$ at $T_{\text{eq}} \sim (m_{3/2}^2 M_{\text{P}}^2 / m_{\tilde{g}})^{1/3} \sim 10^{12}$ GeV. The decrease of the gauge coupling also crucially affects the production of gravitinos after inflation [6] which we now discuss.

The thermal production of gravitinos by gluons, gluinos, quarks and squarks is governed by the Boltzmann equations. The collision term has been calculated to leading order in the gauge coupling. For the gauge group $\text{SU}(N)$, with $2n_f$ chiral multiplets in the fundamental representation, one has [7],

$$\begin{aligned} \frac{dn_{3/2}}{dt} + 3Hn_{3/2} &= C_{3/2}(T, \phi) \\ &= \frac{3\zeta(3)}{32\pi^3} g^2(N^2 - 1) \frac{T^6}{M_{\text{P}}^2} \left(1 + \frac{m_{\tilde{g}}^2}{3m_{3/2}^2}\right) \mathcal{F}(T), \end{aligned} \quad (27)$$

where

$$\begin{aligned} \mathcal{F}(T) &= \left(\ln \left(\frac{T^2}{m_{\text{gluon}}^2(T)} \right) + 0.3224 \right) (N + n_f) \\ &\quad + 0.5781n_f, \end{aligned} \quad (28)$$

with the thermal gluon mass

$$m_{\text{gluon}}^2(T) = \frac{g^2}{6} (N + n_f) T^2. \quad (29)$$

For the gauge coupling we use $g(T, \phi_T)$, except in case of the gluon mass which enters only logarithmically.

In the supersymmetric standard model gravitino production is dominated by QCD, the strong interactions, where we have $N = 3$ and $n_f = 6$. If the gravitino is the LSP and the GUT relations for gaugino masses hold, one has $m_{3/2} \ll m_{\tilde{g}}$. Integrating Eq. (27) up to a reheating temperature $T_R > T_*$, assuming a power decrease of the gauge coupling,

$$g^2(T, \phi_T) \simeq \frac{g_0^2(T)}{1 + (T/T_*)^\alpha}, \quad (30)$$

one obtains a number density to entropy density ratio of gravitinos which is independent of T_R ,

$$\left. \frac{n_{3/2}}{s} \right|_{T_0} = \frac{C_{3/2}(T_*, 0)}{s(T_*)H(T_*)} I_{(\alpha)}. \quad (31)$$

Here T_0 is the present temperature, $s = (2\pi^2/45) \times g_*(T)T^3$ is the entropy density, with $g_*(T_*) = 915/4$ in the supersymmetric standard model, and

$$I_{(\alpha)} = \int_0^\infty \frac{dz}{(1+z^\alpha)^3} = 0.50\text{--}0.73, \quad (32)$$

for $\alpha = 1, \dots, 4$. Inserting the expression for the collision term in Eq. (27) one finds for the energy density to entropy density ratio of gravitinos ($T_R > T_*$),

$$\begin{aligned} \left. \frac{\rho_{3/2}}{s} \right|_{T_0} &= \frac{135\sqrt{10}\zeta(3)}{64\pi^6} \frac{N^2 - 1}{g_*^{3/2}(T_*)} \frac{T_* m_{\tilde{g}}^2(T_*)}{M_{\text{P}} m_{3/2}} \\ &\quad \times I_{(\alpha)} g_0^2(T_*) \mathcal{F}(T_*). \end{aligned} \quad (33)$$

At temperatures T_R much larger than T_* also contributions involving Yukawa interactions may become important, which remains to be studied.

One can now insert the relation (20) between the temperature T_* and gluino and gravitino masses into Eq. (33), which yields the result ($T_R > T_*$),

$$\begin{aligned} \left. \frac{\rho_{3/2}}{s} \right|_{T_0} &= \frac{135\sqrt{5}\zeta(3)}{64\pi^6} \frac{N^2 - 1}{g_*^{3/2}(T_*)} \\ &\quad \times \left(\frac{m_{\tilde{g}}^{3/2}(\mu)}{M_{\text{P}}^{1/2}} \right) \left(\frac{\xi}{a_2 \eta^2} \right)^{1/4} I_{(\alpha)} \hat{\mathcal{F}}(T_*). \end{aligned} \quad (34)$$

Here we have used the gluino mass at a scale μ as parameter, and $\hat{\mathcal{F}}(T_*) = \mathcal{F}(T_*) g_0^{9/2}(T_*) / g_0^3(\mu)$ is a factor $\mathcal{O}(1)$ which takes gauge couplings and their

running into account. Remarkably, in $\rho_{3/2}/s$ the dependence on the gravitino mass has dropped out. For the dominant QCD contribution $N = 3$ and $a_2 = 3/8$ (cf. (14)). Dividing by the critical density $\rho_{\text{crit}}/s = 3.65h^2 \times 10^{-9} \text{ GeV}$ [8] one finally obtains ($T_R > T_*$),

$$\Omega_{3/2}h^2 = 0.1 \times \left(\frac{m_{\tilde{g}}(1 \text{ TeV})}{1.0 \text{ TeV}} \right)^{3/2} \times \left(\frac{\xi}{\eta^2} \right)^{1/4} I_{(\alpha)} \hat{\mathcal{F}}(T_*). \quad (35)$$

For gaugino mediation one has $\xi/\eta^2 = \mathcal{O}(1)$; in the temperature range $T_* = 10^4\text{--}10^{12} \text{ GeV}$ we estimate $I_{(\alpha)} \hat{\mathcal{F}}(T_*) = 0.5\text{--}2$. It is then very astonishing how close the obtained value for $\Omega_{3/2}h^2$ is to the observed one for cold dark matter for gluino masses $\mathcal{O}(1 \text{ TeV})$. The WMAP Collaboration recently obtained (2σ error), $\Omega_{\text{CDM}}h^2 = (\Omega_m - \Omega_b)h^2 = 0.113^{+0.016}_{-0.018}$ [9]. The relic gravitino density $\Omega_{3/2}h^2$ is shown in Fig. 1 as function of the reheating temperature T_R for different values of $m_{\tilde{g}}$ and $m_{3/2}$. At $T_R \simeq T_*$ the density reaches a plateau whose value is essentially independent of T_R and $m_{3/2}$. The figure clearly shows the scaling $T_* \propto m_{3/2}/\sqrt{m_{\tilde{g}}}$.

One may also use Eq. (35) to determine the range of gluino masses consistent with the WMAP result for cold dark matter. Varying $\Omega_{\text{CDM}}h^2$ and $I_{(\alpha)} \hat{\mathcal{F}}(T_*)$ in

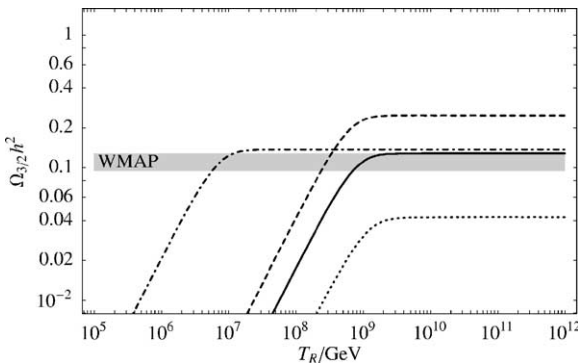


Fig. 1. Relic gravitino density $\Omega_{3/2}h^2$ as function of the reheating temperature T_R for different gravitino and gluino masses: $m_{3/2} = 20 \text{ GeV}$ with $m_{\tilde{g}} = 1.5 \text{ TeV}$ (dashed line), $m_{\tilde{g}} = 1.0 \text{ TeV}$ (full line), $m_{\tilde{g}} = 0.5 \text{ TeV}$ (dotted line), and $m_{3/2} = 200 \text{ MeV}$ with $m_{\tilde{g}} = 1.0 \text{ TeV}$ (dashed-dotted line); $\xi/\eta^2 = 1$, $\alpha = 2$. $\Omega_{3/2}h^2$ reaches a plateau at $T_R \simeq T_* \propto m_{3/2}/\sqrt{m_{\tilde{g}}}$. The band denotes the WMAP result for cold dark matter with a 2σ error.

the ranges specified above we find,

$$m_{\tilde{g}} = (0.5\text{--}2.0) \text{ TeV} \left(\frac{\eta^2}{\xi} \right)^{1/6}. \quad (36)$$

Hence, the hypothesis that gravitinos are the dominant component of dark matter will be tested at LHC!

The range for the gluino mass given in Eq. (36) has been obtained in the case of gaugino mediation where $m_{3/2} = (2\eta/g_0^2)(M/M_P)m_{\tilde{g}}$ (cf. (5) and (6)), with $\eta = \mathcal{O}(1)$. For gravity mediation [10], one obtains the same results with η replaced by $\eta' = \eta M_P/M$. $m_{3/2}$ and $m_{\tilde{g}}$ now have the same order of magnitude, but the gravitino can be the LSP without fine tuning. The range for the gluino mass remains unchanged unless M is smaller than M_P by several orders of magnitude. In the case of gauge mediation [11], η has to be replaced by $\eta' = \eta 8\pi^2 \langle X \rangle / M$ where $\langle X \rangle$ is the messenger scale. The mass range (36) for the gluino mass is then obtained if M is of order $8\pi^2 \langle X \rangle$. Note that the rapid decrease of gauge couplings at high temperature occurs independently of the supersymmetry breaking mechanism.

Our results have important consequences for leptogenesis [12] where the typical baryogenesis temperature is $T_B = \mathcal{O}(10^{10} \text{ GeV})$ or larger [13]. According to previous studies this implies that unstable gravitinos have to be heavier than a few TeV [14,15]. Stable gravitinos may have masses below $\mathcal{O}(1 \text{ keV})$ [16] so that their mass density is below the critical density even when they are thermalized. Further, it has been shown that also gravitino masses $m_{3/2} \sim 10\text{--}100 \text{ GeV}$ can be consistent, which then constrains masses and couplings of other neutralinos and sleptons [17,18].

Our analysis shows that there is no constraint on the reheating temperature for gluino masses below $\mathcal{O}(1 \text{ TeV})$ (cf. Fig. 2) if the gravitino is the lightest supersymmetric particle. For $m_{\tilde{g}} = \mathcal{O}(1 \text{ TeV})$ and reheating temperatures $T_R > T_*$ we find $\Omega_{3/2}h^2 \simeq \Omega_{\text{CDM}}h^2 \simeq 0.1$, independently of $m_{3/2}$.

The maximal value of the critical temperature T_* is obtained for $M \sim M_P$ and $m_{3/2} \sim 100 \text{ GeV}$, so that the gravitino can still be the LSP for a gluino mass $\mathcal{O}(1 \text{ TeV})$. This yields $T_*^{\text{max}} \sim 10^{10} \text{ GeV}$, which happens to coincide with the typical leptogenesis temperature. Hence, for a reheating temperature T_R larger than the leptogenesis temperature T_B , relic gravitinos always have the observed dark matter energy density $\Omega_{\text{CDM}}h^2$ if the gluino mass is $\mathcal{O}(1 \text{ TeV})$. In this way

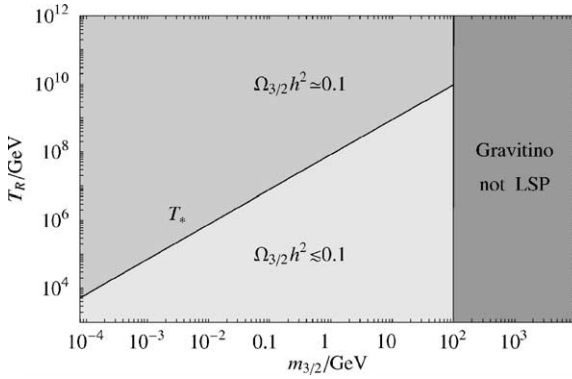


Fig. 2. Relic gravitino density for different values of reheating temperature and gravitino mass. $\xi/\eta^2 = 1$, $m_{\tilde{g}} = 1$ TeV, which implies $m_{3/2} < 0.1$ TeV for a stable gravitino. For $T_R > T_*$, $\Omega_{3/2}h^2$ is independent of T_R and $m_{3/2}$.

the supersymmetry breaking scale in the observable sector is directly determined by the dark matter density $\Omega_{\text{CDM}}h^2$, independently of the supersymmetry breaking scale in the hidden sector!

The interplay of particle physics and cosmology relates some properties of the universe to properties of elementary particles. Of particular interest is the composition of the present energy density Ωh^2 . Leptogenesis explains the baryon density $\Omega_b h^2$ in terms of neutrino masses and mixings. As we have seen, for stable gravitinos the dark matter density $\Omega_{\text{CDM}}h^2$ is then determined by the gluino mass, i.e., the supersymmetry breaking scale in the observable sector, which may also be responsible for the dominant contribution to Ωh^2 , the cosmological constant.

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