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Motivated by the absence of signals of new physics in searches for both new particles at the LHC and a weakly interacting massive particle dark matter candidate, we consider a scenario where supersymmetry is broken at a scale above the reheating temperature. The low-energy particle content then only consists of Standard Model states and a gravitino. We investigate the possibility that the latter provides the main component of dark matter through the annihilation of thermalized Standard Model particles. We focus on the case where its production through scattering in the thermal plasma is well approximated by the nonlinear supersymmetric effective Lagrangian of the associated Goldstino and identify the parameter space allowed by the cosmological constraints, allowing the possibility of a large reheating temperature compatible with leptogenesis scenarios and alleviating the so-called "gravitino problem".

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I. INTRODUCTION

Among the possible hidden symmetries of nature, supersymmetry is one of the most appealing. It provides a candidate for a fundamental theory of nature with a better ultraviolet behavior. If it is realized at low energy, it allows to address the hierarchy problem of the electroweak symmetry breaking sector, provides a dark matter candidate when R parity is conserved, and allows for the unification of the gauge couplings. However, the benefits of this aesthetically attractive scenario need to be questioned in light of increasing tensions with experimental data. On the one hand, there is no sign for new physics in the searches at the LHC, implying strong constraints on the parameter spaces of the models. On the other hand, the negative results in direct and indirect searches for weakly interacting massive particles (WIMPs) are closing the window of parameters corresponding to a neutralino dark matter. This motivates us to push further up the scale of supersymmetry breaking, extending the energy range of validity of the Standard Model. Here, we consider this possibility within a peculiar cosmological scenario.

WIMPs [1] and freeze-in massive particles (FIMPs) [2] are different theoretical frameworks that have been postulated for the production mechanisms of dark matter. Whereas WIMP dark matter are in equilibrium with other particles in the early Universe, when the temperature drops below the dark matter mass they freeze out and form the present relic abundance. In supersymmetric frameworks, the lightest neutralino with a mass around the weak scale is a natural WIMP candidate [1]. The recent analysis on dark matter detections—from indirect methods using its annihilation products (FERMI [3], HESS [4], and AMS [5]) to direct-detection methods using the measurements of nuclear recoils (LUX [6], PANDAX [7], and XENON100 [8])—still did not find any significative signals. However, the sensitivity reached by the different groups begins to exclude large parts of the parameter space predicted by simple WIMPly extensions of the Standard Model (Higgs portal [9], Z portal [10], and even Z' portal [11]). At the supersymmetric (SUSY) level, the authors of Ref. [12] showed that the well-tempered neutralino—which was the most robust supersymmetric candidate in recent years—is now severely constrained by the last results of LUX [6]. A wino-like neutralino is also severely constrained by the latest indirect-detection searches released by the Fermi Collaboration [3]. A nice review of the status of dark matter SUSY searches in different supersymmetric scenarios with neutralino dark matter can be found in Ref. [13]. In any case, the recent prospects exposed by the LUX [14] and FERMI [15] collaborations showed that the WIMP paradigm should be excluded (or discovered) for dark matter masses below 10 TeV in the present generation of detectors.

The lack of experimental signals motivates investigations of other production mechanisms with weaker couplings. FIMP dark matter is an alternative that uses couplings between the dark matter and Standard Model particles suppressed by a much higher scale than the weak scale. Thus they are not in equilibrium with other particles in the early Universe and never reach equilibrium among themselves; their yields keep growing as the temperature decreases. One can distinguish three cases of FIMPs: (i) decay

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products of some heavier particles in equilibrium, (ii) final products of infrared (IR) [$M_{\rm DM}$] dominated processes, and (iii) final products of ultraviolet (UV) [$T_{\rm RH}$] dominated ones. Gravitinos produced by the decay of next-to-lightest supersymmetric particles (NLSPs) [16] belong to the first class. Because the production is dominated at low temperatures, the final yields are largely independent of the reheating temperature. This scenario has a constraint from big bang nucleosynthesis (BBN) since the late decay of NLSPs influences the nucleosynthesis mechanisms that are strongly constrained by the observed abundance of D and ^{4}He in the Universe [17]. The IR-dominated productions usually correspond to renormalizable operators or $2 \to 1$ processes [2]. These are most efficient when the temperature is near the FIMP mass. Thus the yields are not dependent on the reheating temperature. On the other hand, nonrenormalizable operators usually lead to UV-dominated production and the final results are highly dependent on the reheating temperature.

A gravitino is a universal prediction of local supersymmetry models. Its role in cosmology depends on its abundance and its lifetime. Even if in some nonminimal scenarios it can be nonthermally produced at the end of inflation during preheating due to a fine-tuned coupling to the inflaton, the amount expected is model dependent and can be small [18,19]. However, we will not consider this possibility in this work and will instead focus on the case of thermal production. The gravitinos are produced by the scattering of Standard Model states in the thermal plasma after reheating or through the decay of the NLSP. Under the assumption that the reheating temperature is lower than the mass of all of the supersymmetric particles, we are left with only the former possibility. However, the standard scenario of gravitino dark matter suffers from different cosmological difficulties which are referred to collectively as the "cosmological gravitino problems": i) the late decaying superpartners can strongly affect big bang nucleosynthesis [17] and, ii) if thermalized, the relic gravitinos produced overclose the Universe for $m_{3/2} \gtrsim 1$ keV [20], making it difficult to be a warm dark matter candidate if one also takes into consideration large-scale structure formations, the Tremaine-Gunn bound, or Lyman α constraints [21]. We will show that in the high scale supersymmetry framework we propose, where the gravitino is directly produced from the thermal bath scattering, these two issues do not hold anymore. As a consequence, our analysis leads naturally to the prediction of a possibly large reheating temperature $T_{\rm RH}$, usually favored by inflationary or leptogenesis scenarios.

The paper is organized as follows. We establish the framework of our model in Sec. II, insisting on the fundamental mass scales entering in the analysis, before building the effective Lagrangian and computing the cosmological observables in Sec. III. We then conclude in Sec. IV.

II. SUPERSYMMETRY BREAKING AND THE REHEATING TEMPERATURE

In this section we review the different scales relevant for our analysis and their origins: the SUSY-breaking scale, the soft mass terms, the messenger scale, and the gravitino mass.

(1) *The supersymmetry breaking parameters:*

We denote by F the order parameter for supersymmetry breaking, which is a generic combination of auxiliary F or D terms' vacuum expectation values. It corresponds to a spontaneous breaking, and thus implies the existence of a Goldstone fermion: the Goldstino G. The super-Higgs mechanism at work leads to a mass for the gravitino whose value at present time reads [22]

$$m_{3/2} = \frac{F}{\sqrt{3} M_{\rm Pl}}, \tag{1}$$

where $M_{\rm Pl}$ is the reduced Planck mass. The breaking is mediated to the visible sector through messengers lying at a scale $\Lambda_{\rm mess}$. This leads to soft terms of order $M_{\rm SUSY}$:

$$M_{\rm SUSY} = \frac{F}{\Lambda_{\rm mess}}. \tag{2}$$

We shall assume for simplicity that all the masses of sparticles, squarks, sleptons, gauginos, and Higgsinos as well as all the new scalars in the extended Higgs sector are at least of the order of the scale of supersymmetry breaking $M_{\rm SUSY}$. These particles are thus decoupled at reheating time $T_{\rm RH}$. Below $M_{\rm SUSY}$, the particle content is the Standard Model (SM) (with possibly right-handed neutrinos) and the Goldstino. How realistic this assumption is in explicit supersymmetry-breaking models is a model-dependent question. In O'Rafeartaigh models of supersymmetry breaking, the partner of the Goldstino—the sgoldstino $\tilde{G}$—is usually massless at tree level. Quantum corrections are however expected to fix its mass to be one-loop suppressed with respect to the supersymmetry-breaking scale $m^2_{\tilde{G}} \sim \frac{g^2}{16\pi^2} F$, which has to be above the reheating temperature for our model to be self-consistent. In string effective supergravities it is also often the case that the sgoldstino is light, with mass of the order of the gravitino mass [23]. This is also however a model-dependent statement; this can be avoided in models with a large Riemann curvature in the Kahler space [24]. On the other hand, asking for $m_{3/2} \ll M_{\rm SUSY}$ implies

$$\Lambda_{\rm mess} \ll M_{\rm Pl} \tag{3}$$

and in the energy range under consideration the renormalization group equations (RGEs) are those of the SM. In particular, for a Higgs boson of 126 GeV it leads to a vanishing of the quartic coupling at scales of order 2×10^{10} to 3×10^{11} GeV depending on the assumption on the degeneracy of superparticle soft masses, the exact value of the top mass, and the strong interaction gauge coupling (see, for instance, Ref. [25]). We then consider

$$M_{\rm SUSY} \lesssim \{10^{10}\text{–}10^{11}\}\ \text{GeV}. \tag{4}$$

Supersymmetry-breaking scales above this value can be achieved by a modification of the RGEs through the introduction of new light particles. We shall not discuss these cases in details here, as the generalization is straightforward.

(2) *The cosmological parameters:*

The cosmological history described here starts after the Universe is reheated. Some assumptions are made for this epoch: (i) the reheating temperature $T_{\rm RH}$ is small enough to not produce superpartners of the Standard Model particles, and thus $T_{\rm RH}\lesssim M_{\rm SUSY}$; (ii) in the reheating process Goldstinos are scarcely produced. This second condition is a constraint of the nature of the inflaton, its scalar potential, and the branching ratios in its decay. A discussion of the production of Goldstinos at the end of inflation can be found, for example, in Ref. [18].

We consider that the dark matter gravitino interactions are well approximated by the helicity $\pm1/2$ components. This is true in virtue of the equivalence theorem if the energy E of the gravitinos is much bigger than their mass. The enhancement of the interactions of the (longitudinal component of the) light gravitino is a direct consequence of the equivalence theorem between the Goldstino and the longitudinal component of the gravitino $\Psi_\mu \to (\frac{1}{m_{3/2}})\partial_\mu G$, as discussed for the first time in Ref. [26]. Approximating the former by the temperature T of the SM particles in equilibrium leads to the mass hierarchies that define the self-consistency of our setup:

$$m_{3/2} \ll T_{\rm RH} \lesssim M_{\rm SUSY} \lesssim \sqrt{F} \lesssim \Lambda_{\rm mess} \ll M_{\rm Pl}. \tag{5}$$

Note that our bound on the reheating temperature is compatible with thermal leptogenesis. In fact, a lower bound of the reheating temperature is obtained when the latter is identified with the mass of the lightest right-handed neutrino. It is at most of order 10^9 GeV but can be lower depending on the assumptions on the initial abundance and the mass hierarchies of the neutrinos (see, for example, Ref. [27]).

III. GOLDSTINO DARK MATTER

A. Effective Goldstino interactions

Under the assumption $m_{3/2} \ll E \sim T$ discussed above, the gravitino interactions with SM fields are dominated by the helicity $\pm1/2$ components. Moreover, for $E\sim T\lesssim T_{\rm RH}\lesssim M_{\rm SUSY}$, these are described by a nonlinear realization of supersymmetry in the entire observable SM sector, since we will consider all superpartners to be heavy and therefore not accessible in the thermal bath after reheating.[1] The leading-order Goldstino-matter interactions can be divided into two types of contributions: universal [30] and nonuniversal ones [31–33]. We will restrict our analysis to the former, which corresponds to the minimal couplings expected from the low-energy theorem.[2] Their construction starts by defining a "vierbein" [34]

$$e_m^a = \delta_m^a - \frac{i}{2F^2}\partial_m G\sigma^a\bar{G} + \frac{i}{2F^2}G\sigma^a\partial_m\bar{G}, \tag{6}$$

which under a supersymmetry transformation of the parameter ϵ transforms as a diffeomorphism in general relativity,

$$\delta e_m^a = \partial_m\xi^n e_n^a + \xi^n\partial_n e_m^a, \tag{7}$$

where $\xi^n = \frac{i}{F}e_a^n(\epsilon\sigma^a\bar{G} - G\sigma^a\bar{\epsilon})$. The couplings to matter in this original geometrical prescription follow the standard coupling to matter of a metric tensor built out from the vierbein $g_{mn} = \eta_{ab}e_m^a e_n^b$. The corresponding Goldstino-matter effective operators are consequently of dimension eight and take the form

$$L_{2G} = \frac{i}{2F^2}(G\sigma^\mu\partial^\nu\bar{G} - \partial^\nu G\sigma^\mu\bar{G})T_{\mu\nu}, \tag{8}$$

where G is the Goldstino field and $T_{\mu\nu}$ is the energy-momentum tensor of the SM matter fields. The energy-momentum tensor is given by

$$\begin{aligned} T_{\mu\nu} &= +\eta_{\mu\nu}\tilde{L} \\ &\quad + \left[\sum_f\left(-\frac{i}{4}D_\mu\bar{\psi}_f\bar{\sigma}_\nu\psi_f + \frac{i}{4}\bar{\psi}_f\bar{\sigma}_\nu D_\mu\psi_f\right) - D_\mu H D_\nu H^\dagger \right. \\ &\quad \left. + \sum_{\text{SM group}}\frac{1}{2}F_\mu^{a\xi}F_{\nu\xi}^a + (\mu\leftrightarrow\nu)\right]. \end{aligned} \tag{9}$$

The scalar potential and mass terms for scalars and fermions appear in the first term. After the contraction

[1] A reheating temperature below superpartner masses was proposed and investigated in particular in Refs. [28] and [29]. The novelty in our case is that we consider high-scale supersymmetry, so our reheating temperature is much higher compared to these references.

[2] As we will see, our result will not depend drastically on this hypothesis.

between $\eta_{\mu\nu}$ and $G\sigma^\mu\partial^\nu\bar{G}$, the on-shell production of two Goldstinos give a cross section proportional to $m_{3/2}^2$. As $m_{3/2}$ is much smaller than $T_{\rm RH}$, these contributions can be neglected, as we will see later. Then the $2\to 2$ scatterings for the Goldstino production are dominated by the following operators[3]:

$$\frac{i}{2F^2}(G\sigma^\mu\partial^\nu\bar{G}-\partial^\nu G\sigma^\mu\bar{G})(\partial_\mu H\partial_\nu H^\dagger+\partial_\nu H\partial_\mu H^\dagger),$$
$$\frac{1}{8F^2}(G\sigma^\mu\partial^\nu\bar{G}-\partial^\nu G\sigma^\mu\bar{G})$$
$$\times(\bar{\psi}\bar{\sigma}_\nu\partial_\mu\psi+\bar{\psi}\bar{\sigma}_\mu\partial_\nu\psi-\partial_\mu\psi\bar{\sigma}_\nu\psi-\partial_\nu\psi\bar{\sigma}_\mu\psi),$$
$$\times\sum_a\frac{i}{2F^2}(G\sigma^\xi\partial_\mu\bar{G}-\partial_\mu G\sigma^\xi\bar{G})F^{\mu\nu a}F^a_{\nu\xi}, \tag{10}$$

where h, ψ, and $F^a_{\nu\xi}$ stand for a complex scalar (Higgs doublet), gauge bosons, and two-component fermions (quarks and leptons), respectively. Another way to describe the two Goldstinos' interactions with matter is to replace the superpartner soft mass terms with couplings between the Goldstino superfield and the matter superfield multiplets. One can integrate out the heavy (superpartner) components and eliminate them as a function of the light degrees of freedom: the SM fields and the Goldstino. This leads to an effective low-energy theory where the incomplete multiplets are described in terms of constrained superfields [32,35]. The kinetic terms of the sparticles will then lead to dimension-eight operators containing two Goldstinos and two SM fields that generically differ from the ones computed from the low-energy theorem couplings [31]. For the gauge and SM fermion sectors, the resulting cross sections only differ in the angular distribution and numerical constants, whereas the energy dependence is the same as for the low-energy theorem couplings.

Since the masses of the superpartners are of order $M_{\rm SUSY}<\sqrt{F}$, one can worry about effective operators generated after decoupling heavy superpartners with larger coefficients. In particular, there can be dimension-eight operators proportional to $1/M_{\rm SUSY}^4$ and $1/M_{\rm SUSY}^2F$ that would be dominant over the universal couplings we use in our paper. This issue was investigated in the first reference in Ref. [35], where it was shown that by starting from the minimal supersymmetric SM only dimension-eight R-parity-violating couplings of this type are generated. The reason for this is the following: integrating out heavy superpartners (without R-parity violation) leads to factors of $1/M_{\rm SUSY}$ from the propagators (or the square of them for scalar superpartners). On the other hand, the leading interactions of Goldstinos with matter through the soft terms are proportional to $M_{\rm SUSY}/F$ (or the square of them for scalar superpartners). As a result, the factors of $M_{\rm SUSY}$ cancel out, leaving generically dimension-eight operators suppressed by $1/F^2$. The effect of the R-parity-violating couplings on gravitino production was investigated more recently in Ref. [29].

[3]See the Appendix for the expression of these operators in four-component Dirac spinor and γ-matrix notation.

B. Computation of the gravitino relic density

1. The framework

Contrarily to the weakly interacting neutralino, the gravitino falls into the category of feebly interacting dark matter. Its interactions at high energies are governed by the helicity-1/2 component whose couplings are naturally suppressed by the supersymmetry-breaking scale. In gravity-mediated supersymmetry breaking the gravitino is often heavier than the supersymmetric spectrum that it generates. As a consequence, if the gravitino is not sufficiently heavy (i.e., below 30 TeV) it is a long-lived particle which usually decays around the BBN epoch. This gives rise to the famous "gravitino problem" [36,37]. In that case, in order to minimize the observable effects, the gravitino density has to be small enough at the cost of an upper bound on the reheating temperature of the Universe (see, e.g., Ref. [38]). On the other hand, if the gravitino is the LSP it can be a very good dark matter candidate, as either a stable or metastable particle, with a lifetime much longer than the age of the Universe.

The gravitino was in fact the first supersymmetric dark matter candidate ever proposed[4] by Pagels and Primack [20] and Khlopov and Linde [40]. However, they showed that if it is thermalized its mass is restricted to a window of $m_{3/2}\lesssim 1$ keV which places it in the hot scenario, which is in strong tension with current large-scale formation constraints [41] or the Tremaine-Gunn bound ($m_{3/2}\gtrsim 400$ eV) [21]. Then, the authors of Ref. [38] famously showed that the overabundance problem can be avoided if, instead of thermalizing, the gravitino is produced through scattering of a gaugino with a reheating temperature below a critical value, depending on the gaugino spectrum. They obtained

$$\Omega_{3/2}h^2\sim 0.3\left(\frac{1\ \text{GeV}}{m_{3/2}}\right)\left(\frac{T_{\rm RH}}{10^{10}\ \text{GeV}}\right)\sum_i c_i\left(\frac{M_i}{100\ \text{GeV}}\right)^2 \tag{11}$$

where c_i are coefficients of order one, and M_i are the three gaugino masses. We clearly see from Eq. (11) that the density is settled by the reheating temperature. Lower limits on M_3 obtained by the nonobservation of the gluino at the LHC set (for a given gravitino mass) an *upper* limit on the reheating temperature to avoid overclosure of the Universe. These constraints are usually in tension with baryogenesis mechanisms [27], even though some interesting scenarios

[4]To be exact, Fayet [39] already proposed such a hypothesis.

with low reheating temperatures ($T_{\rm RH} \lesssim 100$ TeV) can be found in Ref. [42].

Moreover, later on in Ref. [16] it was shown that another contribution—called the "gravitino freeze-in"—plays an important role. It corresponds to the decay of the superpartners while they are still in thermal equilibrium. Indeed, for a sufficiently large supersymmetric spectrum (which seems to be the case at the LHC) the short lifetime of squarks or sleptons induces this process. The only way to circumvent the overabundance is to lower the reheating temperature *below* the supersymmetric spectrum to deal with the queue of the distribution. However, the origin of the gravitino is still the supersymmetric partners, through their decay. A nice summary can be found in Ref. [43]. Adding the BBN constraints gives an upper bound on the gravitino mass of about 10 GeV [44–46].

All the scenarios discussed above made the hypothesis of thermal production of the gravitino, through supersymmetric partners in thermal equilibrium with the primordial plasma. And, this thermalization hypothesis is the deep source of tension between the cosmological observables (density of dark matter, structure formation, BBN, or leptogenesis) and the data. However, if for some reason the supersymmetric-breaking scale is above the reheating temperature *while still keeping it at a high scale*, the SM superpartners will be too heavy to reach thermal equilibrium. That naturally solves the preceding tension, but the issue of gravitino production remains. Note also that for $T_{\rm RH} \lesssim M_{\rm SUSY}$ thermal corrections to the effective potential may become small enough to lead to only negligible displacements of the scalars' vacuum expectation values from their late-time values. This kind of high-scale SUSY scenario can arise easily and naturally in string-inspired constructions (see, for instance, Ref. [47] and references therein). A way to populate the Universe with gravitinos is through a *direct* freeze-in from the thermal bath itself. In this new scenario, the gravitinos are produced at a rate smaller than the one corresponding to the expansion of the Universe, and therefore they do not have time to reach thermal equilibrium. It "freezes in" in the process of reaching it, as the strong suppression of the scattering cross sections by the scale F^2 in Eq. (10) prevents the gravitinos from being in thermal equilibrium with the Standard Model bath. We propose using this scenario to confront with cosmological data.

2. Gravitino production through freeze-in

From the interaction generated through the Lagrangian (10), one can compute the production rate $R = n_{\rm eq}^2 \langle \sigma v \rangle$ of the gravitino $\tilde{G}$, generated by the annihilation of the Standard Model bath of density $n_{\rm eq}$. The details of the computation are developed in the Appendix [Eq. (A7)], and we obtain

$$R = \sum_i n_{\rm eq}^2 \langle \sigma v \rangle_i \simeq 21.65 \times \frac{T^{12}}{F^4}. \qquad (12)$$

The Boltzmann equation for the gravitino density $n_{3/2}$ can be written as

$$\frac{dY_{3/2}}{dx} = \left(\frac{45}{g_* \pi}\right)^{3/2} \frac{1}{4\pi^2} \frac{M_P}{m_{3/2}^5} x^4 R, \qquad (13)$$

with $x = m_{3/2}/T$, $Y_{3/2} = n_{3/2}/\mathbf{s}$, where $\mathbf{s}$ is the entropy density, and g_* is the effective number of degrees of freedom thermalized at the time of gravitino decoupling (106.75 for the Standard Model). Here, we use the Planck mass $M_P = 1.2 \times 10^{19}$ GeV. We then obtain after integration

$$Y_{3/2} = \frac{21.65 M_{\rm P} T_{\rm RH}^7}{28\pi^2 F^4} \left(\frac{45}{g_* \pi}\right)^{3/2} \simeq 3.85 \times 10^{-3} \frac{M_{\rm P} T_{\rm RH}^7}{F^4}. \qquad (14)$$

The relic abundance

$$\Omega h^2 = \frac{\rho_{3/2}}{\rho_c^0} = \frac{Y_{3/2} s_0 m_{3/2}}{\rho_c^0} \simeq 5.84 \times 10^8 Y_{3/2} \left(\frac{m_{3/2}}{1\ \text{GeV}}\right) \qquad (15)$$

is then

$$\Omega_{3/2} h^2 \simeq 0.11 \left(\frac{100\ \text{GeV}}{m_{3/2}}\right)^3 \left(\frac{T_{\rm RH}}{5.4 \times 10^7\ \text{GeV}}\right)^7. \qquad (16)$$

As can be seen, the dependence on the reheating temperature is completely different from the case where the gravitino is produced through the scattering of the gaugino in Eq. (11). A similar behavior can be observed in the SO (10) framework [48] or in extended neutrino sectors [28]. In all of these models the production process appears at the beginning of the thermal history, and is then very mildly dependent on the hypothesis or the physics appearing *after* reheating. The reheating temperature is then a prediction of the model (for a given gravitino mass) once one applies the experimental constraints of WMAP [49] and Planck [50]. Another interesting point is that a look at Eqs. (14) and (16) shows that even the dependance on the particle content is very mild. Indeed, due to the large power $T_{\rm RH}^7$, the total number of degrees of freedom (or even channels) does not influence the final reheating temperature that much, which is predicted to be around 10^8 GeV for a gravitino with the electroweak scale. Even the hypothesis of universal couplings [30] or nonuniversal ones [31,32] will not drastically affect our Eq. (16).

Our result is plotted in Fig. 1 in the plane ($m_{3/2}$, $T_{\rm RH}$) [in the plane ($T_{\rm RH}$, $M_{\rm NLSP}$) in Fig. 2] where we represent the parameter space allowed by the relic abundance constraints $\Omega_{3/2} h^2 \simeq 0.12$ [49,50]. As can be seen, a large part of the parameter space is allowed by cosmology, giving reasonable values of $T_{\rm RH} \simeq 10^5$–10^{10} GeV for a large range of gravitino masses (MeV–PeV). The region below the orange (dashed) line is excluded as the gravitino would be too heavy to be produced by the freeze-in mechanism, whereas the region above the green (dotted) line corresponds to a freeze-out scenario. In the latter region, the production cross section $\langle\sigma v\rangle$ is sufficiently high to reach thermal equilibrium. This occurs when $n\langle\sigma v\rangle \gtrsim H(T_{\rm RH}) \simeq T_{\rm RH}^2/M_{\rm Pl}$. A quick look at Eq. (12) shows that such a large cross section is obtained for a high reheating temperature or small values of F (and thus a light gravitino), explaining the shape of the green region in Fig. 1. However, once the gravitino is in thermal equilibrium, its density is given by the classical freeze-out (FO) mechanism

$$\Omega_{3/2}^{\rm FO} = \frac{n_{3/2} m_{3/2}}{\rho_c^0} \Rightarrow \simeq 0.1 \left(\frac{m_{3/2}}{180\ \text{eV}}\right), \qquad (17)$$

which corresponds obviously to the intersecting point in Fig. 1.

There exists another potential nonthermal source of gravitino production: the decay of the NLSP. Indeed, this contribution also exists in the standard supersymmetric framework, through the relic abundance produced by the decay of the NLSP (usually a sfermion $\tilde{f}$) into $\tilde{f} \to Gf$. This process (being proportional to $n_{\tilde{f}}^{\rm eq}$) is highly Boltzmann suppressed in our scenario where $T_{\rm RH} \ll M_{\rm NLSP}$, but there still exists some parameter space where the NLSP is in equilibrium. Then the production of Goldstinos is a combination of the decay of the NLSP, QCD processes, and the SM freeze-in. An analysis in this scenario with a very low $T_{\rm RH}$ ($\lesssim$GeV) can be found in Ref. [29].

C. Comments on the *R*-parity violation operators

R-parity violation operators can also be introduced in the high-scale supersymmetry scenario discussed in this work.

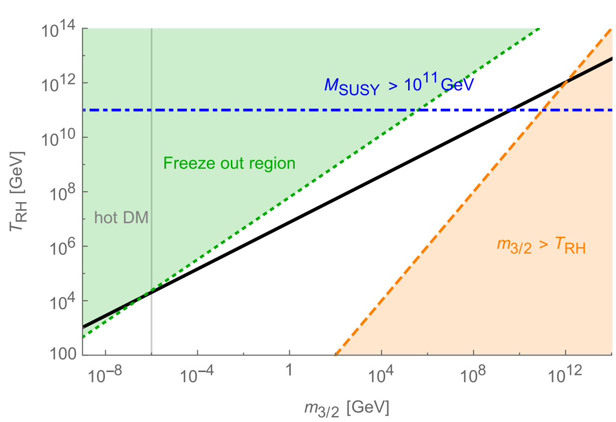

FIG. 1. Region in the parameter space ($m_{3/2}$; $T_{\rm RH}$) respecting the relic abundance constraint [49,50] from Eq. (16). The points above the black line are excluded because the gravitino would overclose the Universe. The blue line shows the constraint from the Higgs mass with an observed value of 125 GeV, which sets an upper limit for the scale of supersymmetry breaking [Eq. (4)].

The corresponding operators involving Goldstino fields include dimension-five operators as

$$\frac{\mu_i}{F} l_I^i \sigma^\mu \bar{G} D^\mu h^j + \text{H.c.}, \qquad (18)$$

dimension-six ones as [51]

$$\frac{iC^I}{F} \epsilon_{ij} (l_I^i \partial_\mu G) D^\mu h^j + \text{H.c.}, \qquad (19)$$

and dimension-eight operators of the form

$$\frac{\lambda''_{ijk}}{m_i^2 F} u_i d_j \Box(d_k G), \quad \frac{\lambda'_{ijk}}{m_i^2 F} q_i l_j \Box(d_k G), \quad \frac{\lambda_{ijk}}{m_i^2 F} l_i l_j \Box(e_k G), \qquad (20)$$

plus permutations. Here μ_i and C are dimensionful and dimensionless coefficients, respectively, l_I are the three lepton doublets in the SM, and m_i^2 are soft terms of the heavy superpartners that were integrated out. The $2 \to 1$ gravitino production through these operators will be suppressed at temperatures higher than the gravitino mass and only become important at late times; therefore, they do not need to be considered for the production of gravitino dark matter.

When R parity is violated, the gravitinos are no longer stable but can decay, giving rise to observable signatures. The latter are independent of the production mechanisms and the previous analyses in the literature apply to our case. The relevant operators can be derived from the above but should be written using the gravitino field. Since the heavy supersymmetric particles decouple in our case, the coefficients of the R-parity-violating operators are not necessarily constrained from preserving baryon asymmetry, as in previous studies [52].

However, one characteristic of our construction is that it allows for a very heavy gravitino (above the PeV scale).

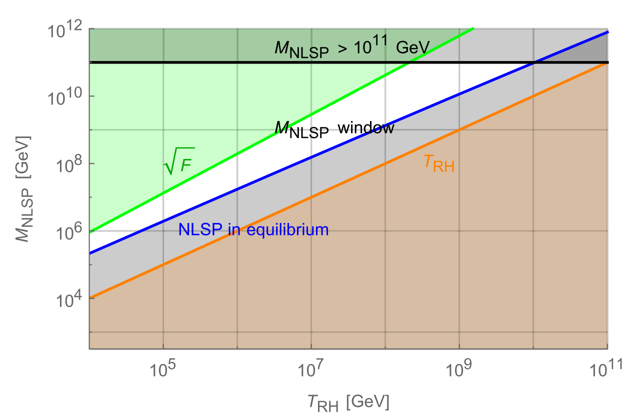

FIG. 2. The parameter space for $M_{\rm NLSP}$. It must be lower than $\sqrt{F}$. The blue line corresponds to $n_g \langle\sigma v\rangle_{gg\to\tilde{g}\tilde{g}} = H$. Below the line, the NLSP is still in equilibrium and can decay to a gravitino.

Smoking-gun signals like $G \to h\nu$ and $G \to \gamma\nu$ can be observable with telescopes like IceCube for neutrinos [53] or the future Cerenkov Telescope Array for the photon [54]. In both cases, a monochromatic high-energy signal should be the signature of gravitino decay, the spatial morphology distinguishing decaying dark matter (proportional to its density ρ) from annihilating dark matter (proportional to ρ^2).

IV. CONCLUSION

We considered the framework of high-scale supersymmetry, where the scale of superpartners $M_{\rm SUSY}$ lies above the reheating temperature, whereas the gravitino mass $m_{3/2}$ stays below it. In this case, there still exist processes which thermally produce gravitinos through scattering of the Standard Model particles at the earliest time of reheating. Our result is well summarized by Fig. 1 and Eq. (16), where one can observe and understand the strong dependence of the relic abundance on the reheating temperature $T_{\rm RH}$. Our result predicts a large reheating temperature ($\sim 10^8$ GeV for a ~ 100 GeV gravitino). This scale pattern $m_{3/2} \ll T_{\rm RH} \ll M_{\rm SUSY}$ is common in some string models with high-scale supersymmetry breaking [47] and opens new possibilities in model building.

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APPENDIX

1. Computing the gravitino production rate *R*

In this appendix we provide the details of the computation of the annihilation rate $n_{\rm eq}^2\langle\sigma v\rangle$. Indeed, after symmetrization and by switching to four-component fermionic notation, one can extract from Eq. (10) the effective Lagrangian

$$\begin{aligned}\mathcal{L} \supset &-\frac{i}{2F^2}\left(\partial_\mu\bar{G}\gamma_\nu\frac{1+\gamma_5}{2}G - \bar{G}\gamma_\nu\frac{1+\gamma_5}{2}\partial_\mu G\right)\\ &\times(\partial^\mu H^\dagger\partial^\nu H + \partial^\nu H^\dagger\partial^\mu H)\\ &+\frac{1}{8F^2}\left(\partial_\mu\bar{G}\gamma_\nu\frac{1+\gamma_5}{2}G - \bar{G}\gamma_\nu\frac{1+\gamma_5}{2}\partial_\mu G\right)\\ &\times\left(\bar{\Psi}\gamma^\nu\frac{1+\gamma_5}{2}\partial^\mu\Psi - \partial^\mu\bar{\Psi}\gamma^\nu\frac{1+\gamma_5}{2}\Psi\right.\\ &\left.+\bar{\Psi}\gamma^\mu\frac{1+\gamma_5}{2}\partial^\nu\Psi - \partial^\nu\bar{\Psi}\gamma^\mu\frac{1+\gamma_5}{2}\Psi\right)\\ &-\frac{i}{2F^2}\left(\partial_\mu\bar{G}\gamma_\nu\frac{1+\gamma_5}{2}G - \bar{G}\gamma_\nu\frac{1+\gamma_5}{2}\partial_\mu G\right)F^{\mu\lambda a}F^{\nu a}_{\lambda},\end{aligned} \tag{A1}$$

where G is the Goldstino in a four-component Dirac fermion notation, and H, Ψ, and $F_{\mu\nu}$ are the Higgs field, Standard Model fermions, and gauge field strength, respectively. It is then straightforward to compute the averaged production rate R for the process $1+2\to3+4$ in the case of early decoupling, when all the particles i in the thermal bath, of temperature T, are relativistic ($m_i \ll T \Rightarrow E_i = p_i$):

$$R_i = n_{\rm eq}^2\langle\sigma v\rangle_i = \int f_1 f_2 d\cos\beta\frac{E_1E_2dE_1dE_2}{1024\pi^6}\int|\mathcal{M}|_i^2 d\Omega,$$

where $f_i = \frac{1}{e^{E_i/T}\pm1}$ for a fermionic (bosonic) distribution, β is the angle between the colliding particles 1 and 2 with energies E_1 and E_2, respectively, in the laboratory frame, and Ω is the solid angle between the incoming particle 1 and outgoing particle 3 in the center of mass frame.[5] From Eq. (A1) one can easily deduce

$$|\bar{\mathcal{M}}|_h^2 = \frac{s^4}{16F^4}(\cos^2\theta - \cos^4\theta), \tag{A2}$$

$$|\bar{\mathcal{M}}|_f^2 = \frac{s^4}{256F^4}(1+cos\theta)^2(1-2\cos\theta)^2, \tag{A3}$$

[5] See Refs. [55] and [2] for details.

$$|\bar{\mathcal{M}}|^2_V = \frac{s^4}{128F^4}(2-\cos^2\theta-\cos^4\theta) \tag{A4}$$

for the scalar, fermionic, and vectorial contributions, respectively.[6] The *total* averaged production rate $n^2_{\rm eq}\langle\sigma v\rangle$ is then given by

$$R = \sum_i n^2_{\rm eq}\langle\sigma v\rangle_i = 4n^2_{\rm eq}\langle\sigma v\rangle_h + 45n^2_{\rm eq}\langle\sigma v\rangle_f + 12n^2_{\rm eq}\langle\sigma v\rangle_V, \tag{A5}$$

$$\begin{aligned} n^2_{\rm eq}\langle\sigma v\rangle_h &= \frac{48\zeta(6)^2}{\pi^5F^4}T^{12} = \frac{48\pi^7}{(945)^2F^4}T^{12}, \\ n^2_{\rm eq}\langle\sigma v\rangle_f &= \frac{72\zeta(6)^2}{\pi^5F^4}\left(\frac{31}{32}\right)^2T^{12} = \frac{72\pi^7}{(945)^2F^4}\left(\frac{31}{32}\right)^2T^{12}, \\ n^2_{\rm eq}\langle\sigma v\rangle_V &= \frac{264\zeta(6)^2}{\pi^5F^4}T^{12} = \frac{264\pi^7}{(945)^2F^4}T^{12}, \end{aligned} \tag{A6}$$

implying

$$R = \frac{6400[\zeta(6)]^2}{\pi^5F^4}T^{12} = \frac{6400\pi^7}{(945)^2F^4}T^{12} \simeq 21.65\times\frac{T^{12}}{F^4}. \tag{A7}$$

[6]The integration on the phase space should be treated with care, as the lorentz invariant $s=(P_1+P_2)^2=2P_1.P_2=2E_1E_2(1-\cos\beta)$ in the laboratory frame and $\int\frac{x^n}{e^x-1}=n!\zeta(n+1)$.

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