Decay rates of unstable particles and the extreme energy cosmic rays top-down scenarios

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We provide a unified formula for the quantum decay rate of heavy objects (particles) whatever they may be: topological and nontopological solitons, *X* particles, cosmic defects, microscopic black holes, fundamental strings, as well as the particle decays in the standard model. Extreme energy cosmic ray (EECR) top-down scenarios are based on relics from the early Universe. The key point in the top-down scenarios is the necessity to adjust the lifetime of the heavy object to the age of the Universe. This *ad hoc* requirement needs a very high dimensional operator to govern its decay and/or an extremely small coupling constant. The arguments produced to fine-tune the relic lifetime to the age of the Universe are critically analyzed. The natural lifetimes of such heavy objects are, however, microscopic times associated with the grand unified theory energy scale ($\sim 10^{-28}$ sec or shorter). It is at this energy scale (by the end of inflation) that they could have been abundantly formed in the early Universe, and it seems natural that they decayed shortly after being formed. The annihilation scenario for EECRs (''wimpzillas'') is also considered and its inconsistencies analyzed.

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I. INTRODUCTION

We provide a unified description for the quantum decay formula of unstable particles which encompass topological and nontopological solitons, *X* particles, cosmic defects, microscopic black holes, and fundamental strings, as well as the particle decays in the standard model (muons, Higgs bosons, etc.). In all cases the decay rate can be written as

$$\Gamma = \frac{g^2 m}{\text{numerical factor}}, \tag{1.1}$$

where g is the coupling constant, m is the typical mass in the theory (it could be the mass of the unstable particle), and the numerical factor often contains relevant mass ratios for the decay process.

Top-down scenarios for extreme energy cosmic rays (EECRs) are based on heavy relics from the early Universe, which are assumed to decay at the present time, or on topological defects also originating in the early Universe. For all relics (whatever their nature: heavy particles, topological and nontopological solitons, black holes, microscopic fundamental strings, cosmic defects, etc.), one has to fine-tune the lifetime of these objects to be the age of the Universe.

The second type of top-down scenario relies on the existence of a network of topological defects formed during phase transitions in the early Universe. Such topological defects should survive until the present to produce the observed EECRs. If they decay in the early Universe we go back to the previous case. It must first be noted that only some grand unified field theories support topological defects. Moreover, recent cosmic microwave background (CMB) anisotropy measurements from the Boomerang, Maxima, Dasi, Archeops, and Wilkinson Anisotropy Probe Experiments (WMAP) [1] have seen no evidence of topological defects, strongly disfavoring their eventual presence in the present Universe.

A third type of top-down scenario proposes that stable heavy relics can produce EECRs through annihilation by pairs [2,3]. This scenario suffers from a different type of inconsistency as we show below.

Heavy relics could have been formed by the end of inflation at typical grand unified theory (GUT) energy scales, but their natural lifetime would be of the order of microscopic times typically associated with the GUT's energy scales [4]. Further top-down scenarios are reviewed in Ref. [5].

The problem in the top-down scenarios is not the formation of heavy particles or topological defects. They could all have been generated in the early Universe. The key problem is their existence today (i.e., their imposed lifetime of the order of the age of the Universe) and the value of their mass, which must be adjusted to be $\sim 10^{20}$ eV.

The key drawback of the first top-down scenarios is the lifetime problem. The *ad hoc* requirement of a lifetime of the order the age of the Universe for the heavy particles implies an operator with a very high dimension describing the decay, and/or an extremely small coupling constant (see Sec. VII). As discussed in Sec. VII, a number of so-called ''solutions'' have been invoked in order to cope with the lifetime problem, but all these ''remedies'' replace one assumption by another one.

The second type (topological defect networks) and third type (annihilation of heavy relics) of EECR top-down scenarios suffer from equally severe drawbacks, as discussed in Sec. VIII.

A common feature of top-down approaches is that the arguments trying to support a long lifetime for the *X* particles successively call for more and more speculative explanations.

EECRs may result from the acceleration of protons and ions by shock waves in astrophysical plasmas (Fermi acceleration mechanism) [6]. That is, charged particles can be efficiently accelerated by electric fields in astrophysical shock waves [6–9]. This is the so-called diffusive shock acceleration mechanism, yielding a power spectrum with

$$n(E)\sim E^{-\alpha}, \tag{1.2}$$

with $2.3\leqslant\alpha\leqslant 2.5$–$2.7$. This spectrum is well verified over 13 orders of magnitude in energy.

Large enough sources or sources with strong enough magnetic fields can accelerate particles to the energies of the observed EECRs. Sources in the vicinity of our galaxy such as hot spots of radio galaxies (working surfaces of jets and the intergalactic medium) [10,11], the fireballs producing gamma ray bursts (GRBs) [12], magnetars (young neutron stars) with strong magnetic fields [13], and blazars (active galactic nuclei with relativistic jets directed along the line of sight) like BL Lacertae can evade the Greisen-Zatsepin-Kuz'min (GZK) bound.

Extreme energy cosmic rays have been observed by a number of experiments at energies above 10^{20} eV [14]. Forthcoming cosmic rays detectors like the Auger array and the EUSO and OWL space observatories are expected to greatly improve our present knowledge of the EECRs gathered from Fly's Eye, HiRes, AGASA, and previous detectors [14,15].

II. TOPOLOGICAL SOLITONS, NONTOPOLOGICAL SOLITONS, AND HEAVY PARTICLES

Stable solutions in classical field theory (such as monopoles) become (heavy) particles in quantum field theory. There is no difference at the quantum level between heavy particles associated with a local field and those associated with classical stable solutions.

The stability of classical solutions in field theory is a highly nontrivial issue. There are basically two types of solution: topological and nontopological. Topological classical solutions are associated with a nonzero topological number (topological charge) which vanishes for the vacuum. If there is a lower bound for the energy of the solution involving this topological number, the classical solution is stable. This is the case for kinks in one-space-dimension scalar theories, vortices in the two-dimensional Higgs model [16], monopoles in the three-dimensional Georgi-Glashow model [17], and Hopf solitons in appropriate three-dimensional scalar models [18]. In all known cases, classical stability comes together with quantum stability.

Gravitational analogues of these classical solutions exist in the Euclidean (imaginary time) regime [19,20]: they are black holes in three space dimensions (with periodicity in the imaginary time), which are the gravitational analogues of electric type monopoles, and Taub-Nut solutions in four space dimensions (gravitational analogues of magnetic type monopoles). The topological charges here are related to the temperature and magnetic charge of the solutions, respectively [20,21].

It must be stressed that the mere presence of a conserved topological number does not guarantee the stability of the corresponding classical solution. The energy must be related to the topological number in question such that a nonzero topological number implies a nonzero energy [22]. Otherwise, a classical solution possessing nonzero topological number can decay into lighter particles.

In other words, the topological charge may be disconnected from the dynamics and it can decay in the course of the evolution. A topological soliton may collapse, losing its topological charge. This does not happen when the topological charge bounds the energy from below.

Nontopological solitons are stable thanks to a conserved $U(1)$ charge of "electric" type [23]. Again, the mere presence of a conserved $U(1)$ charge does not guarantee stability for charged particles except for the lightest one. Let us indicate by m the mass of the lightest charged particle and let us take its $U(1)$ charge as the unit of charge. Assume that there are heavier particles with mass $M>m$ and charge $Q>1$ with $M=M(Q)$. A sufficient condition for quantum stability is

$$M(Q)<mQ,$$

since a particle with charge Q and mass larger than or equal to mQ can always decay into Q particles of mass m and unit charge, respecting charge and energy conservation.

It must be stressed that in quantum theory all nonforbidden processes do happen.

III. QUANTUM DECAY OF HEAVY PARTICLES

Typically, the decay of a heavy particle with mass m_X can be described by an effective interaction Lagrangian formed by the local field $X(x)$ associated with this heavy particle times the lighter fields into which it decays. Let us take muon decay, which is a well known case. Notice the mass of the muon $m_\mu=206.8m_e\gg m_e$.

The effective Fermi Lagrangian can be written as [24]

$$\mathcal{L}_I=-\frac{G_F}{\sqrt{2}}\overline{\psi}_{\nu_\mu}\gamma^\alpha(1+\gamma_5)\psi_\mu\overline{\psi}_{\nu_e}\gamma_\alpha(1+\gamma_5)\psi_{\nu_e}, \tag{3.1}$$

where ψ_μ stands for the muon field, and ψ_{ν_e} and ψ_{ν_μ} for the electron neutrino and muon neutrino fields, respectively. The Fermi coupling G_F has the dimension of an inverse square mass.

The muon width Γ_μ describing the decay is then given by

$$\Gamma_\mu=\frac{G_F^2 m_\mu^5}{192\pi^3}.$$

The Fermi coupling can be related to the W mass as follows:

$$\frac{G_F}{\sqrt{2}}=\frac{g^2}{8m_W^2},$$

where g, the standard model coupling, is dimensionless. Thus,

$$\Gamma_\mu=\frac{g^4 m_\mu}{6144\pi^3}\left(\frac{m_\mu}{m_W}\right)^4. \tag{3.2}$$

As we shall see below, Eq. (3.2) has the generic structure of the decay width of an unstable particle.

For muon decay, the monomial interaction in the effective Lagrangian (3.1) has dimension 6 in mass units.

An analogous example is the Higgs boson decay into muons, neutrinos, $W^{\pm}$, and the Z^0. Notice that the Higgs boson mass m_H must be higher than the $W^{\pm}$ mass m_W and the Z^0 mass. The Lagrangian as given by the standard model is here

$$2g \sin\theta_W M_W H W^+_\mu W^\mu_- , \tag{3.3}$$

and a similar expression for the coupling with the Z. Here θ_W stands for Weinberg's angle.

One finds for the Higgs boson decay rate [24,25] into a $W^{\pm}$ pair

$$\Gamma_{Higgs\to W^+W^-}=\frac{g^2}{64\pi}m_H\left(\frac{m_H}{m_W}\right)^2\sqrt{1-4\left(\frac{m_W}{m_H}\right)^2}\times\left[1-4\left(\frac{m_W}{m_H}\right)^2+12\left(\frac{m_W}{m_H}\right)^4\right]. \tag{3.4}$$

If $m_H \gg m_W$ we can simplify this formula to

$$\Gamma_{Higgs\to W^+W^-}=\frac{g^2}{64\pi}m_H\left(\frac{m_H}{m_W}\right)^2. \tag{3.5}$$

We here consider the limit of a heavy Higgs boson for illustration, although it is known by now that the Higgs boson is not a lot heavier than the W.

The monomial interaction in the effective Lagrangian (3.3) has dimension 3 in mass units.

Notice that in both cases, Eq. (3.2) and Eq. (3.4), the width grows as a positive power of the mass of the decaying particle.

Let us consider an effective Lagrangian containing a local monomial of dimension n (in mass units)

$$\mathcal{L}_I=\frac{g}{M^{n-4}}X\Theta. \tag{3.6}$$

Here the field X is associated with the decaying particle of mass m_X and Θ stands for the product of fields coupled to it.

Then, the decay rate for a particle of mass m_X takes the form

$$\Gamma=\frac{g^2}{\text{numerical factor}}m_X\left(\frac{m_X}{M}\right)^{|2n-8|}. \tag{3.7}$$

Γ_μ [Eq. (3.2)] and Γ_{Higgs} [Eq. (3.5)] (if we assume $m_H \gg m_W$) correspond to $n=6$ and $n=3$, respectively. When $m_X \sim M$, the right hand side of Eq. (3.7) can be multiplied by a function of M/m_X as in Eq. (3.4).

IV. QUANTUM DECAY OF SOLITONS

The mass of classical soliton solutions (like magnetic monopoles in unified theories) is of the form

$$M_{sol}=\frac{\mu}{g^2}$$

where μ is the mass of the basic fields in the Lagrangian and g their dimensionless coupling. For small coupling these objects are much heavier than the particles associated with the basic fields in the Lagrangian.

Quantum mechanically, the soliton mass acquires corrections of order g^0 and higher. To one-loop level one finds

$$M_{sol}=\frac{\mu}{g^2}+\frac{1}{2}\sum_n[\omega_n-\omega_n^0], \tag{4.1}$$

where ω_n stands for the frequency of oscillations around the soliton. These oscillations are close but not identical to the frequency of oscillations around the vacuum ω_n^0. The sum in Eq. (4.1) yields a finite result proportional to μ [22].

Now, if the classical solution is unstable, some of the frequencies ω_n develop an imaginary part $i\mu\beta$ where β is a pure number. Hence,

$$\text{Im } M_{sol}=\beta\mu \quad \text{and} \quad \text{Re } M_{sol}=\frac{\mu}{g^2}+\mathcal{O}(g^0) \tag{4.2}$$

and

$$\Gamma_{sol}=\text{Im } M_{sol}=g^2\beta \text{ Re } M_{sol}. \tag{4.3}$$

We see that the width Γ_{sol} has a similar structure as for heavy particles in the previous section.

The term $\mathcal{O}(g^0)$ in Eq. (4.2) stands for the first quantum correction to the mass. Notice that we choose $\hbar=1$, which is absorbed in g^2.

V. QUANTUM DECAY OF FUNDAMENTAL STRINGS

The decay of closed strings in string theory has been computed to the dominant order (one string loop) [26]. Assuming the closed string in an Nth excited state, it can decay into lower excited states including the graviton, the dilaton, and the massless antisymmetric tensor. The mass of this quantum string is given by

$$m^2=32\pi TN,$$

where T is the string tension $T=1/(4\pi\alpha')$ and α' the string constant. The length of such string is given by $L=2\alpha' m$ (see, however, Refs. [27]). Here, we consider only closed bosonic strings. Analogous formulas hold for the total decay rate of open bosonic strings and for superstrings. The corresponding decay products being different in each case, i.e., for superstrings, the appropriate superposition of decay products includes gluinos, axions, and quarks.

One finds, for the total width of string decay [26],

$$\Gamma_{string}=\frac{\kappa^2\sqrt{TN}}{\text{numericalfactor}}, \tag{5.1}$$

where the dimensionless coupling κ is given by

$$\kappa=48\pi\sqrt{2GT}.$$

The total width can then be rewritten as

$$\Gamma_{string}=\frac{\kappa^2 m}{1083\times\text{numerical factor}}. \tag{5.2}$$

This formula again has the same structure as the previous widths (3.7) and (4.3) once we identify $g=\kappa, m_X=\text{Re }M_S=m$.

Equation (5.2) can be rewritten as

$$\Gamma_{string}=42\frac{GTm}{\text{numerical factor}}=\frac{21}{16\pi}\frac{Gm^3}{\text{numerical factor}}$$
$$=\frac{21\sqrt{2}}{\text{numerical factor}}\frac{\sqrt{N}G}{\alpha'^{3/2}}. \tag{5.3}$$

VI. QUANTUM DECAY OF BLACK HOLES

As is known, in the context of field theory black holes decay semiclassically through thermal emission at the Hawking temperature [28]:

$$T_{BH}=\frac{\hbar c}{4\pi k_B}\frac{1}{R_s}, \quad R_s=\frac{2GM}{c^2}$$

(M being the black hole mass and k_B the Boltzmann constant).

Black hole emission follows a "graybody" spectrum (the "filter" being the black hole absorption cross section $\sim R_s^2$). The mass loss rate in this process can be estimated following a Stefan-Boltzmann relation,

$$\frac{dM}{dt}=-\sigma R_s^2 T_{BH}^4\sim T_{BH}^2,$$

where σ is a constant. Thus, the black hole decay rate is

$$\Gamma_{BH}=\left|\frac{1}{M}\frac{dM}{dt}\right|\sim GT_{BH}^3\sim\frac{G}{R_s^3}.$$

As evaporation proceeds, the black hole temperature increases until it reaches the string temperature [29]

$$T_{string}=\frac{\hbar c}{k_B}\frac{1}{bL_s}, \quad L_s=\sqrt{\frac{\hbar\alpha'}{c}}$$

(L_s being the fundamental string length and b a constant exclusively depending on the spacetime dimensionality and the string model chosen). The black hole enters its string regime $T_{BH}\to T_{string}$, $R_s\to L_s$, becomes a string state and decays with the width

$$\Gamma_{BH}\to GT_s^3\sim\frac{G}{\alpha'^{3/2}}\sim\Gamma_{string}.$$

Notice that this formula is similar to Eq. (5.3) and again has the generic structure of the widths Eqs. (3.7)–(4.3) and (5.2) if one identifies $g=\kappa$, $m_X=\text{Re }M_S=m$.

We consider here both fundamental strings and black holes since their decay rates can be nicely recast as in Eq. (1.1) independently of whether or not they may be considered as candidate sources of EECRs.

VII. LIFETIMES OF PARTICLES AND THE AGE OF THE UNIVERSE

Heavy particles with masses in the GUT scale can be produced in large numbers during inflation and just after inflation [30]. The production mechanism is parametric or spinodal amplification in the inflaton field; that is, linear resonance of the quantum modes of the heavy field in the background or condensate of the inflaton. In addition, nonlinear quantum phenomena play a crucial role and can enhance the particle production [4]. Such nonlinear production is of the same order of magnitude as the gravitational production of particles by the time dependent metric.

Once these heavy particles are produced, they must have a lifetime of the order of the age of the Universe in order to survive in the present Universe and decay into ultrahigh energy cosmic rays. Only in the early Universe is the production of such heavy objects feasible due to their large mass.

Moreover, in order to be the source of EECRs, these particles must have the mass of the observed EECRs, namely, $m_X>10^{21}$ eV$=10^{12}$ GeV.

Let us assume that the effective Lagrangian (3.6) describes the decay of the X particles [31]. Their lifetime will be given by Eq. (3.7):

$$\tau_X=\frac{\text{numerical factor}}{g^2}\frac{1}{m_X}\left(\frac{M}{m_X}\right)^{2n-8}$$
$$=\frac{\text{numerical factor}}{g^2}\frac{1}{m_X}10^{6(n-4)},$$

where we set the GUT mass $M=10^{15}$ GeV. The age of the Universe is $\tau_{Universe}\sim 2\times 10^{10}$ yr and we have to require that $\tau_X>\tau_{Universe}$. Therefore,

$$10^{54}<\frac{\text{numerical factor}}{g^2}10^{6(n-4)} \quad \text{or} \quad \log_{10} g<3(n-13), \tag{7.1}$$

and we dropped the numerical factor in the last step.

For $g\sim 1$, Eq. (7.1) requires an operator Θ with dimension at least 13 in the effective Lagrangian (3.6), which is a pretty high dimension. That is, one needs to exclude all operators of dimension lower than 13 in order to extremely suppress the decay. Clearly, one may accept lower dimension operators Θ, paying the price of a small coupling g. For example, $g=10^{-9}$ and $n=10$ satisfy the above bound of a rather high dimension operator. Notice that a moderate n such as $n=4$ lowers the coupling to $g\sim 10^{-27}$.

In summary, a heavy X particle can survive from the early Universe until the present time if one chooses an extremely small coupling g and/or an operator Θ with high enough dimension.

None of these assumptions can be supported by arguments other than imposing a lifetime of the age of the Universe on the X particle. That is, the lifetime must here be fine-tuned: one has to build an *ad hoc* Lagrangian to describe the X-particle decay. Indeed, a variety of *ad hoc* Lagrangians have been proposed in the literature together with symmetries that can adjust a wide variety of lifetimes [32].

In order to cope with the lifetime problem a number of so-called ''solutions'' have been invoked, but all the ''remedies'' replace one assumption by another, as follows.

(1) An assumed new global symmetry to protect the X particle and which would be broken nonperturbatively only by quantum gravity wormholes [33] or instanton-type effects [34] to make τ_X = age of the Universe. However, these quantum gravity effects are poorly controlled, and basic theoretical uncertainties remain (the sign of the Euclidean gravity action being only one of them), which would produce the opposite effect to the one claimed; thus τ_X would be exponentially shortened (instead of increased) by wormholes.

(2) Discrete gauge symmetries to construct high dimensional operators [35]. As stated before, these are all *ad hoc* Lagrangians built on the assumption that such group symmetries could have a physical role. No fundamental physical reason exists to argue for them.

(3) Along the same line of thinking, fractionally charged particles (''cryptons'') have been invoked [36] from some particularly chosen hidden sectors of particular effective string/M theory inspired models. Then, support for the assumption of a long lifetime for τ_X would come from a strongly interacting (bound state) sector and its nonpertubative dynamics (which is not controlled), in flipped $SU(5)$, for instance. But, as for many ''particle'' sectors appearing in string inspired phenomenology, no physical reason exists to choose such states, in particular fractionally charged (bound) state particles.

Finally, along the line of reasoning of (1)–(3) above, comparison with the stability or ''metastability'' of the proton has been invoked in order to support the stability (and decay) of the X particles. In the standard model, in which baryon-lepton number is conserved, proton decay can be realized only by introducing *ad hoc* nonrenormalizable high dimensional operators; then, a new global symmetry is invoked for the X particle, broken only by operators suppressed by M^n with $M = M_{Planck}$ and $n > 7$ [37].

In other words, *ad hoc* proton decay is argued to support *ad hoc X* particle decay. As is known, GUT models predict proton decay (which has not been found so far), placing a lower bound on the proton lifetime of $\tau_{proton} > 1.6\times10^{25}$ yr. Proton decay is, however, a natural consequence of grand unification as leptons and quarks belong to the same multiplet.

A common feature of the top-down approaches is that the arguments trying to support a long lifetime for the X particles successively call for more and more speculative explanations. Still, the essential question in the top-down scenarios remains, i.e., if X particles and topological defects in such scenarios did not decay in the early Universe, shortly after they formed, why should they decay just now?, i.e., the lifetime fine-tuning remains.

The top-down scenarios are just tailored to explain the observed events. There is absolutely no physical reason to assume that relics have such a mass (and not any other value) and such a lifetime.

VIII. COSMIC DEFECTS AND HEAVY PARTICLES

The question of the stability of topological solutions is a highly nontrivial issue. The mere presence of a conserved topological charge does *not* guarantee their stability; the energy must be related to the topological charge and must be bounded from below by the topological charge. Otherwise, the topological defect is unstable.

Closed vortices from Abelian and non-Abelian gauge theories are not topologically stable in $3+1$ spacetime dimensions. Static vortices in $3+1$ spacetime dimensions just collapse to a point since their energy is proportional to their length. They do that in a very short (microscopic) time.

It must be noticed that only a restricted set of spontaneously broken non-Abelian gauge theories exhibit vortex solutions. For example, there are no topologically stable vortices in the standard $G = SU(3)\times SU(2)\times U(1)$ model in $3+1$ spacetime dimensions just because $\Pi_1(G)$ and $\Pi_2(G)$ are trivial for such group manifolds. (For a recent review, see [38].) Grand unified theories may or may not possess vortex solutions in $2+1$ spacetime dimensions depending on which representations of the gauge group contain the Higgs fields.

Cosmic strings are closed vortices of horizon size. In $3+1$ spacetime dimensions, strings collapse very fast except if they are of horizon size in which case their lifetime is of the order of the age of the Universe. However, such horizon size cosmic strings are excluded by the CMB anisotropy observations and by the isotropy of cosmic rays.

Such gigantic objects behave classically whereas microscopic closed strings (for energies $< M_{Planck} = 10^{19}$ GeV) behave quantum mechanically.

The existence of cosmic string networks is not established although they have been the subject of much work. If such networks existed in the early Universe they may have produced heavy particles X of the type discussed before, and all the discussion of their lifetime applies here. The discussion of the lifetime problem also applies to rotating superconducting strings, which have been proposed as classically stable objects [39].

In summary, a key point here is the instability of topological defects in $3+1$ spacetime dimensions. Unless one chooses very specific models [17–20,23], topological defects decay even classically with a short lifetime. They collapse to a point at a speed of the order of the speed of light in $3+1$ spacetime dimensions.

A. Annihilation top-down scenarios

There are top-down scenarios where, instead of decay, annihilation of the relic superheavy particles (wimpzillas) has been proposed [2,3]. That is, in this scenario the relics with mass $M_X \sim 10^{12}$ GeV are stable and produce EECRs through annihilation when they collide [2,3]. Here, the lifetime free parameter is replaced by the annihilation cross section. These superheavy particles are assumed to be produced during reheating [2]. Its annihilation cross section is thus bounded by the amount of dark matter in the Universe: $\sigma_X \sim \alpha(M_X)^{-2}$ and $\alpha \leqslant 0.01$. In Ref. [3] the EECR flux produced is computed for several scenarios.

For a smooth dark matter distribution assuming a Navarro-Frenk-White (NFW) singular profile, it is required that $\sigma_X = 6\times10^{-27}$ cm^2 in order to reproduce the observed EECR flux [3]. But this value for σ_X is 10^{27} larger than the maximum value compatible with $\Omega \sim 1$ in the early Universe (corresponding above to $\alpha \sim 0.01$).

In a second scenario, dark matter is assumed to form into clumps with NFW profiles. In this way, the EECR flux is about ~ 2000 times larger due to the increase of the dark matter density [3]. However, that needs again a value for σ_X 10^{24} times larger than its value in the early Universe.

Finally, isothermal clumps have been considered. There, the EECR flux turns out to be independent of σ_X but the internal radius of the clump R_{min} is proportional to σ_X. For our galaxy, R_{min} turns out to be $\sim 10^{-42}$ times the size of the clump. For a 10 kpc clump this gives $R_{min} \sim 10^{-20}$ cm, which is a high energy quantum microscopic scale.

A further problem in this scenario raised in Ref. [3] is that the predicted EECR flux is too high by a factor of 10^{15}. Then, in order to reduce the flux, it is proposed in Ref. [3] that these wimpzillas amount to only 10^{-15} of the dark matter. This can be achieved by setting $\alpha \sim 10^{-17}$ in the annihilation cross section. But such σ_X makes a problem for R_{min} which then becomes too small at $\sim 10^{-37}$ cm$\sim 10^{-4} \times$ Planck length.

B. Signatures of top-down scenarios

We have discussed above the main points of view from which the top-down approach can be theoretically criticized. Let us now mention some characteristic features of top-down models which can be taken as signatures to constrain or disallow them from observational data.

(1) The spectra of the particles generated in the top-down models are typically flatter than in the bottom-up ones. In contrast to the acceleration mechanisms, the top-down generated spectra do not follow a power law.

(2) The composition of the EECR's at the source in the top-down scenarios is dominated by gamma rays and neutrinos (only 5% of the energy is in protons). Although propagation over cosmological distances modifies the ratio of gamma rays to protons, photons still considerably dominate over protons. Top-down scenarios could be constrained by the cascade produced at low energies (MeV–GeV) by gamma rays originating at distances larger than the absorption length [40].

(3) The fluxes of EECR's provided by topological defect models are much lower than required. Simulations of cosmic string networks including self-intersection, intercommutation, multiple loop fragmentation, as well as cusp annihilation, all produce fluxes too low as compared with observations [40,41]. [Some simulations of long string networks claiming flux enhancement were recently discussed, but since the typical distance between two such string segments is of the Hubble scale, the EECR's produced in this way would be completely absorbed [40,41]. And if, by chance, a string as such were near us (about a few tens of Mpc), it would imply a large anisotropy in EECR events, which is not observed.] In any case, simulations of the dynamics of cosmic string networks (with the cosmic expansion included) are not well controlled, making them not predictive.

The present data from the HiRes and AGASA experiments seem incompatible with each other above 10^{20} eV, where AGASA has eight events beyond the GZK bound. Except for these AGASA events, the available EECR spectrum today seems compatible with the GZK effect showing up as predicted [6,8].

Let us recall that the recent CMB anisotropy observations [1] strongly disfavor topological defects in the present Universe. (We have discussed topological defects here since they are still considered in the EECR top-down literature.)

IX. CONCLUSIONS

In summary, if the X particles, whatever their origin and type, could be made sufficiently stable to survive until now, then their decay products could provide the EECRs observed today. However, an X particle lifetime of the order of the age of the Universe must be imposed *ad hoc* i.e., fine-tuned, while the natural lifetime for those particles should be extremely short, about 10^{-28} sec at most.

Various GUTs contain candidates for X particles of masses around the GUT scale ranging approximately from 10^{12} GeV to 10^{16} GeV depending on the model. These particles could have been produced naturally in the early Universe, typically by the end of inflation [30]. Analogously, topological defects, fundamental strings, and primordial black holes could have been formed in the early Universe. The hard job, however, is to have these heavy objects still present and decaying today. Instead of that, it seems more natural that the X particles and the other heavy objects mentioned above decayed in the early Universe shortly after being formed, having lifetimes corresponding to their respective energy scales. Their decay products then form relic primordial backgrounds, such as graviton, neutrino, and dilaton backgrounds, as we now know the relic photon CMB background. Those backgrounds could have characteristic detectable spectra and signatures containing information about the early Universe.

Finally, the standard model of cosmic ray acceleration (diffusive shock acceleration) based on Fermi ideas explains the nonthermal power energy spectrum of CRs over at least 13 orders of magnitude. It is reasonable to extend such mechanisms to EECRs and this seems plausible. However, stimulating physical and astronomical problems remain in understanding and explaining the CR spectrum well below extreme energies.

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