Helium photodisintegration and nucleosynthesis: Implications for topological defects, high energy cosmic rays, and massive black holes

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We consider the production of ^{3}He and ^{2}H by ^{4}He photodisintegration initiated by nonthermal energy releases during early cosmic epochs. We find that this process cannot be the predominant source of primordial ^{2}H since it would result in anomalously high ^{3}He/D ratios in conflict with standard chemical evolution assumptions. We apply this fact to constrain topological defect models of highest energy cosmic ray (HECR) production. Such models have been proposed as possible sources of ultrahigh energy particles and γ rays with energies above 10^{20} eV. The constraints on these models derived from ^{4}He photodisintegration are compared to corresponding limits from spectral distortions of the cosmic microwave background radiation and from the observed diffuse γ-ray background. It is shown that for reasonable primary particle injection spectra and within the models considered previously superconducting cosmic strings, unlike ordinary strings or annihilating monopoles, cannot produce the HECR flux at the present epoch without violating at least the ^{4}He-photodisintegration bound. The constraint from the diffuse γ-ray background rules out the dominant production of HECR by the decay of grand unification particles in models with cosmological evolution assuming standard fragmentation functions. Constraints on massive black hole induced photodisintegration are also discussed.

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I. INTRODUCTION

In this paper we consider various constraints inferred from the possible photodisintegration of ^{4}He in the early universe. Following Protheroe, Stanev, and Berezinsky [1] we note that the photodisintegration of this isotope can be employed to place stringent limits on early cosmic energy injections associated with, for example, decaying particles [2,3], evaporating black holes [4], or annihilating topological defects [5–10]. Our focus here will be particularly on constraining the latter scenario. It has also been suggested that ^{4}He photodisintegration in the early universe could be a production mechanism for the observed light-element abundances of deuterium and ^{3}He [11]. In this work we will study the feasibility of such a scenario and show that the (^{3}He/^{2}H) ratio poses a problem to it. We will show that photodisintegration yields (^{3}He/^{2}H)$\gg 1$ and since ^{2}H is destroyed and ^{3}He increases with evolution, measures of (^{3}He/^{2}H) place severe constraints on photodisintegration.

Nonthermal energy releases at high redshifts may leave various observable signatures. The cosmic microwave background radiation (CMBR) has been measured to have a blackbody spectrum to very high accuracy [12]. Any injection of energy between redshifts of $z \simeq 10^3$ and $z \simeq 3 \times 10^6$ may produce observable spectral distortions of the blackbody spectrum [13]. Here the lower redshift represents the approximate epoch of decoupling (assuming no reionization), whereas the higher redshift represents the epoch at which double-Compton scattering is still efficient enough to completely thermalize significant energy releases [14].

The diffuse γ-ray background observed at the present epoch can also be used to constrain early cosmic energy injections [15]. For redshifts $z \lesssim$ 300–1000 pair production by γ rays on protons and ^{4}He is rare so that the universe becomes transparent to γ rays with energies below $E_{\rm max}$. Here the energy $E_{\rm max}$ is

$$E_{\rm max} \simeq \frac{m_e^2}{15T} \simeq 17 \text{ GeV} \left(\frac{T}{1 \text{ eV}}\right)^{-1} , \qquad (1)$$

where T is the CMBR temperature and m_e is the electron mass. $E_{\rm max}$ is related to the threshold energy for e^+e^--pair creation by high-energy γ rays scattering off CMBR photons. Any radiation with energies above this threshold is effectively instantaneously "recycled" by pair production ($\gamma\gamma_{\rm CMBR} \to e^+e^-$) and inverse Compton scattering of the created electrons and positrons ($e\gamma_{\rm CMBR} \to e\gamma$). These processes yield a degraded γ-ray spectrum with generic energy dependence $\propto E_\gamma^{-1.5}$ considerably below $E_{\rm max}$ before steepening and finally cutting off at $E_{\rm max}$ [3]. Significant energy releases in form of high-energy γ rays and charged particles at epochs with redshifts below $z \simeq$ 300–1000 may therefore produce a present day γ-ray background and are subject to constraint.

For redshifts smaller than $z \simeq 10^6$ stringent constraints on various forms of injected energy can also be derived from the possible photodisintegration of ^{4}He and the concomitant production of deuterium and ^{3}He. The injec-

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tion of high-energy particles and γ rays above the energy threshold $E_{\rm max}$ will initiate an epoch of cascade nucleosynthesis subsequent to the epoch of standard primordial nucleosynthesis at $T \sim 100$ keV. The abundance yields of ^{2}H and ^{3}He produced by ^{4}He photodisintegration during cascade nucleosynthesis are quite independent from the primary γ-ray and charged particle energy spectra. Deuterium and ^{3}He abundance yields depend only on the amount of injected energy and the injection epoch. For the detailed calculations leading to these conclusions the reader is referred to the work by Protheroe, Stanev, and Berezinsky [1]. The nucleosynthesis limits on the release of energy into the primordial gas can be up to a factor of ~ 100 more stringent than equivalent limits on energy releases derived from distortions of the CMBR-blackbody spectrum.

For redshifts $z \gtrsim 10^6$, corresponding to CMBR temperatures of $T \gtrsim 200$ eV, the photodisintegration of ^{4}He is inefficient. This is because the energy threshold for pair production falls below the energy threshold for ^{4}He photodisintegration, $E_{\rm max} \lesssim E_{\rm th}^{^4{\rm He}}$. The best nucleosynthesis limits on decaying particles and annihilating topological defects in the cosmic temperature range 1 keV $\lesssim T \lesssim 10$ keV come from the possible photodisintegration of deuterium [3,16]. These limits are stronger than analogous limits from distortions of the CMBR blackbody spectrum.

In this narrow temperature range limits on decaying particles and topological defects may, in fact, be more stringent due to effects of injecting antinucleons. Antinucleons may be produced during $\gamma\gamma_{\rm CMBR}$ pair production for γ energies $E_\gamma \gtrsim 10^5$ GeV or when there is a significant hadronic decay channel for a massive decaying particle or topological defect. These antinucleons can then annihilate on ^{4}He thereby producing approximately equal amounts of ^{2}H and ^{3}He [17]. We will, however, not further pursue this idea here.

For temperatures above $T \simeq 1$ keV there are virtually no constraints on decaying particles and topological defects from distortions of the CMBR blackbody spectrum. However, stringent limits on decaying particles and topological defects may obtain from the injection of hadrons (for a review see [3]). An injection of mesons and baryons generally increases the neutron-to-proton ratio and results in increased ^{4}He-mass fractions (1 MeV $\gtrsim T \gtrsim 100$ keV) and/or increased ^{2}H and ^{3}He abundances (100 keV $\gtrsim T \gtrsim 10$ keV; [18]). It has been suggested that a combination of ^{4}He hadrodestruction and ^{2}H, ^{3}He photodestruction induced by a late-decaying particle ($T \sim 3$ keV) may bring big-bang-produced light-element abundances close to observationally inferred abundance constraints for a wide range of fractional contributions of baryons to the closure density, Ω_b [19].

The observational signatures of such scenarios are primordial isotope ratios of (^{3}He/^{2}H)$\simeq$ 2–3 and ^{6}Li/^{7}Li~ 1, contrasting the predictions of a standard, or inhomogeneous, big-bang freeze-out from nuclear statistical equilibrium. For a wide range of parameters, such as decaying particle lifetimes and hadronic branching ratios, these models would overproduce ^{2}H and ^{3}He and therefore the calculations by Dimopoulos *et al.* [19] do also serve as constraints on particle parameters and abundances. We note here that the high (^{3}He/^{2}H) ratio may in fact be a severe problem for such scenarios.

In this paper we restrict ourselves to constraints derived from the effects of nonthermal energy injections at epochs with redshifts $z \lesssim 10^6$. The outline of the paper is as follows. In Sec. II we briefly review the observationally inferred light-element abundances of ^{2}H and ^{3}He. We then consider ^{4}He-photodisintegration scenarios and their compatibility with the observations. In Sec. III we study the effects of possible energy injection by superconducting strings, ordinary strings, and magnetic monopoles on the primordial ^{2}H and ^{3}He abundances, the distortions of the CMBR blackbody, and the diffuse γ-ray background. In these scenarios we assume that such topological defects would radiate on a level such that they could produce the observed highest energy cosmic rays at the present epoch. Conclusions are drawn in Sec. IV. Throughout this paper we will mostly use $c = \hbar = 1$.

II. CONSTRAINTS ON ^{4}He PHOTODISINTEGRATION AS THE PREDOMINANT SOURCE OF PRIMORDIAL DEUTERIUM

In this section we investigate scenarios which have ^{4}He photodisintegration as an efficient production mechanism of the light elements deuterium and ^{3}He. In this study we are naturally led to consider the primordial ratio of (^{3}He/^{2}H)$_p$. This is because the ratio of these light isotopes emerging from the big band nucleosynthesis (BBN) process, (^{3}He/^{2}H)$_{\rm BBN}$, is quite different from that emerging from the ^{4}He photodisintegration, (^{3}He/^{2}H)$_{\rm photo}$. In particular, we expect generic isotope ratios of (^{3}He/^{2}H)$_{\rm BBN} \lesssim 1$, and (^{3}He/^{2}H)$_{\rm photo} \gg 1$. We will show that this fact can be used to severely constrain the photodisintegration of ^{4}He as the principal source of primordial deuterium. We will also show that the observationally inferred abundances of ^{2}H and ^{3}He may imply a factor 2–3 more stringent constraints on the primordial number densities of decaying particles and on the energy injected by topological defects than previous work has assumed.

The most accurate determination of a (^{3}He/^{2}H) ratio is thought to come from solar system observations of ^{3}He abundances. Geiss [20] reanalyzed the existing data and inferred for the abundances of deuterium and ^{3}He at the time of solar system formation

$$1.2\times10^{-5} \lesssim \left(\frac{^3{\rm He}}{^1{\rm H}}\right)_\odot \lesssim 1.8\times10^{-5}\ ,$$

$$1.6\times10^{-5} \lesssim \left(\frac{^2{\rm H}}{^1{\rm H}}\right)_\odot \lesssim 3.3\times10^{-5}\ , \qquad (2)$$

$$0.34 \lesssim \left(\frac{^3{\rm He}}{^2{\rm H}}\right)_\odot \lesssim 1.13\ .$$

A determination of the interstellar medium abundances of ^{2}H and ^{3}He is less precise due to observational difficulties [21]. The observed (^{2}H/H) ratios ranges between $5 \times 10^{-6} \lesssim (^2\mathrm{H/H})_{\mathrm{ISM}} \lesssim 2 \times 10^{-5}$ [22]. Interstellar (^{3}He/H) ratios are observed in the range $1.1 \times 10^{-5} \lesssim (^3\mathrm{He/H})_{\mathrm{ISM}} \lesssim 4.5 \times 10^{-5}$ [23]. These abundances imply a present (^{3}He/^{2}H)-isotope ratio of $0.55 \lesssim (^3\mathrm{He}/^2\mathrm{H})_{\mathrm{ISM}} \lesssim 9$.

Deuterium is the most fragile of the light isotopes. It is easily destroyed during the pre-main sequence evolutionary stage of stars via $^2\mathrm{H}(p,\gamma)^3\mathrm{He}$. Furthermore, there are no plausible galactic production sites for deuterium. Epstein, Lattimer, and Schramm [24] summarize the arguments against a galactic origin of deuterium. The chemical evolution of ^{3}He is less clear. It is known that ^{3}He is destroyed to some extent in massive stars ($M \gtrsim 5$–$8 M_\odot$), whereas low-mass stars ($M \lesssim 1$–$2 M_\odot$) may be net producers of ^{3}He. This theory is supported by the observation of ^{3}He abundances in planetary nebulae. It is certainly very reasonable to assume that standard chemical evolution models can only increase the primordial $(^3\mathrm{He}/^2\mathrm{H})_p$ ratio:

$$\left(\frac{^3\mathrm{He}}{^2\mathrm{H}}\right)_t \gtrsim \left(\frac{^3\mathrm{He}}{^2\mathrm{H}}\right)_p . \tag{3}$$

In this expression $(^3\mathrm{He}/^2\mathrm{H})_t$ denotes the isotope ratio at some cosmic time t and the primordial isotope ratio $(^3\mathrm{He}/^2\mathrm{H})_p$ includes any pregalactic production mechanism, such as big bang nucleosynthesis and ^{4}He photodisintegration in the early universe. Note that the inferred (^{3}He/^{2}H) ratios at the time of solar system formation and the present epoch are consistent with the assumption of monotonically increasing (^{3}He/^{2}H) ratios with time.

The (^{3}He/^{2}H) ratio in a standard homogeneous big bang nucleosynthesis (SBBN) scenario at baryon-to-photon ratio $\eta = 3 \times 10^{-10}$ is $(^3\mathrm{He}/^2\mathrm{H})_{\mathrm{SBBN}} \simeq 0.2$. An upper limit on the (^{3}He/^{2}H) ratio in SBBN can be obtained by requiring the ^{4}He-mass fraction to satisfy $Y_p \lesssim 0.25$, whereas a lower limit on this isotope ratio can be estimated from the conservative bound $(^2\mathrm{H/H}) \lesssim 3 \times 10^{-4}$. This yields the SBBN range

$$0.09 \lesssim \left(\frac{^3\mathrm{He}}{^2\mathrm{H}}\right)_{\mathrm{SBBN}} \lesssim 0.55 . \tag{4}$$

Typical (^{3}He/^{2}H)-isotope ratios resulting in homogeneous big bang scenarios are not very different from those in Eq. (4).

The detailed calculations by Protheroe, Stanev, and Berezinsky [1] show that the abundance ratio of (^{3}He/^{2}H) produced during cascade nucleosynthesis in the early universe exceed

$$\left(\frac{^3\mathrm{He}}{^2\mathrm{H}}\right)_{\mathrm{photo}} \gtrsim 8 , \tag{5}$$

for a wide range of fractional contributions of baryons to the closure density, Ω_b, Hubble parameters H_0 in units of 100 km sec^{-1} Mcp^{-1}, h, and epochs of energy injection. This is because in ^{4}He photodisintegration the effective cross sections for the two-nucleon photoabsorption processes [$^4\mathrm{He}(\gamma,pn)^2\mathrm{H}$ and $^4\mathrm{He}(\gamma,^2\mathrm{H})^2\mathrm{H}$] are roughly ten times smaller than the effective cross sections for the single-nucleon photoabsorption processes [$^4\mathrm{He}(\gamma,p)^3\mathrm{H}$ and $^4\mathrm{He}(\gamma,n)^3\mathrm{He}$] [25].

Note that Eq. (5) applies strictly only under the following assumption. In cascade nucleosynthesis it is assumed that the main fraction of radiation is injected above the energy threshold Eq. (1) for $\gamma\gamma_{\mathrm{CMBR}} \to e^-e^+$ pair creation. Pair creation and inverse Compton scattering will then yield a generic γ-ray spectrum with energy dependence $\propto E_\gamma^{-1.5}$ below $E_{\mathrm{max}}/2$ and $\propto E_\gamma^{-5}$ above before cutting off at E_{max}. These γ rays can be effective in photodisintegrating ^{4}He where the competing process is the consumption of γ rays by Bethe-Heitler pair production on hydrogen and helium.

When radiation is injected below E_{max} the γ rays may have a spectrum quite different from the behavior $\propto E_\gamma^{-1.5}$ depending on the actual γ-ray source. In principle, it is then conceivable to photodisintegrate ^{4}He in such a way that isotope ratios of $(^3\mathrm{He}/^2\mathrm{H}) \simeq 1$ result. This could be accomplished by a γ-ray source which preferentially radiates above energies of $E \simeq 100$ MeV but below E_{max}. This is because only in the energy range between the ^{4}He-photodisintegration threshold $E_{\mathrm{th}}^{^4\mathrm{He}} = 19.8$ MeV and $E \simeq 100$ MeV the effective cross section for ^{3}He production in ^{4}He photodisintegration is roughly ten times larger than the effective cross section for ^{2}H production in this process. For γ-ray energies $E \gtrsim 100$ MeV these cross sections are roughly equal. In practice, any such scenario has to occur at relatively low redshifts $z \lesssim 10^3$ so that there will not develop a "softer" second generation γ-ray spectrum produced by Bethe-Heitler pair production and inverse Compton scattering. Note that for redshifts $z \lesssim 10^3$ the universe becomes transparent to γ rays. In this case, however, significant deuterium production would require γ-ray fluxes which would exceed the present day diffuse γ-ray background.

This can be understood by the following rough estimate. When the universe is transparent to γ rays the rate of change in the deuterium abundance due to ^{4}He photodisintegration is

$$(1+z)^3 \frac{d\left(\frac{n_{^2\mathrm{H}}(z)}{(1+z)^3}\right)}{dt} \approx \sigma_\gamma n_{^4\mathrm{He}}(z) 4\pi \int_{E_{\mathrm{th}}}^{\infty} dE(z) j_\gamma(z) , \tag{6}$$

leading to a (^{2}H/^{1}H) ratio at the present epoch

$$\left(\frac{n_{^2\mathrm{H}}}{n_{^1\mathrm{H}}}\right) \approx \frac{4\pi \Delta E(z)}{(1+z)} j_\gamma(z=0) \sigma_\gamma \left(\frac{n_{^4\mathrm{He}}}{n_{^1\mathrm{H}}}\right) (1+z)^{3/2} t_0 . \tag{7}$$

In these expressions n denote number densities for the various nuclear species, $\sigma_\gamma \approx 10^{-1}$ mb is the photodisintegration cross section for the processes [$^4\mathrm{He}(\gamma,pn)^2\mathrm{H}$ and $^4\mathrm{He}(\gamma,^2\mathrm{H})^2\mathrm{H}$], t_0 is the age of the universe, and j_γ

is the photon number flux per unit energy and steradian. In order to produce a number ratio $(^2\mathrm{H}/^1\mathrm{H}) \approx 10^{-5}$ by a hard γ-ray spectrum injected at $(1+z) \approx 10^3$ a photon flux of $j_\gamma(z=0) \sim 10\ \mathrm{MeV}^{-1}\,\mathrm{cm}^{-2}\,\mathrm{s}^{-1}\,\mathrm{sr}^{-1}$ is needed. Here we assumed that $j_\gamma(z \approx 10^3)$ peaks around 1 GeV with $\Delta E(z) \approx 1$ GeV. This should be compared to $j_\gamma(E \approx 1\ \mathrm{MeV}) \sim 10^{-2}\ \mathrm{MeV}^{-1}\,\mathrm{cm}^{-2}\,\mathrm{s}^{-1}\,\mathrm{sr}^{-1}$ [15] (at the redshifted energy $E \approx 1$ MeV) observed at the present epoch.

It should be noted that γ rays could also be effective in photodisintegrating ^{3}He and ^{2}H and thereby in resetting any initial $(^3\mathrm{He}/^2\mathrm{H})_{\mathrm{photo}}$-isotope ratio produced during cascade nucleosynthesis to smaller values. However, the relative abundances of ^{4}He targets to ^{3}He targets is approximately 10^3–10^4 to 1, so that for roughly equal photodisintegration cross sections the number densities of γ rays in the energy range between the ^{3}He-photodisintegration threshold $E_{\mathrm{th}}^{3\mathrm{He}} = 5.4$ MeV and $E_{\mathrm{th}}^{4\mathrm{He}} = 19.8$ MeV should be 10^3–10^4 times larger than the number densities of γ rays with energies above $E_{\mathrm{th}}^{4\mathrm{He}}$. Such a scenario would require an extremely "soft" γ-ray spectrum.

We can derive limits on the allowed contributions of ^{4}He photodisintegration to the primordial ^{2}H and ^{3}He abundances. This can be done by employing the solar system $(^3\mathrm{He}/^2\mathrm{H})$-isotope ratio from Eq. (2) and assuming that this ratio represents a conservative upper limit on the primordial $(^3\mathrm{He}/^2\mathrm{H})$-isotope ratio [refer to Eq. (3)]. Note that when either one of Eqs. (2) or (3) does not apply one of the widely used standard assumptions of galactic chemical evolution has to break down. We can derive an upper limit on the fraction of deuterium $f_{^2\mathrm{H}}^{\mathrm{photo}}$ contributed to the primordial deuterium abundance by ^{4}He photodisintegration. A simple calculation of the abundance average then yields

$$f_{^2\mathrm{H}}^{\mathrm{photo}} \lesssim \frac{\left(\frac{^3\mathrm{He}}{^2\mathrm{H}}\right)_{\odot} - \left(\frac{^3\mathrm{He}}{^2\mathrm{H}}\right)_{\mathrm{BBN}}}{\left(\frac{^3\mathrm{He}}{^2\mathrm{H}}\right)_{\mathrm{photo}} - \left(\frac{^3\mathrm{He}}{^2\mathrm{H}}\right)_{\mathrm{BBN}}} \ . \tag{8}$$

By using the upper limit for $(^3\mathrm{He}/^2\mathrm{H})_\odot$ from Eq. (2), the lower limit in Eq. (4) for the $(^3\mathrm{He}/^2\mathrm{H}\ )_{\mathrm{SBBN}}$ ratio, and Eq. (5) for the $(^3\mathrm{He}/^2\mathrm{H})_{\mathrm{photo}}$ ratio we derive

$$f_{^2\mathrm{H}}^{\mathrm{photo}} \lesssim 13\% \ . \tag{9}$$

It is evident that the contribution of deuterium produced during cascade nucleosynthesis to the total primordial deuterium abundance has to be small in order to not overproduce ^{3}He. This also implies that generic ^{4}He-photodisintegration scenarios *cannot be the predominant production mechanism* of the primordial ^{2}H and ^{3}He light-element abundances. The stringent limit of Eq. (9) can only be evaded if either there existed an extremely "soft" γ-ray source in the early universe or if generic features of the galactic destruction or production of ^{3}He and ^{2}H are for some yet unknown reason not understood.

Gnedin and Ostriker [11] have proposed the interesting scenario of a very early formation ($z \simeq 800$) of massive black holes. If these black holes do accrete material which emits a quasarlike x-ray and γ-ray spectrum they may induce the photodisintegration of ^{4}He and the reionization of the universe. The reionization of the universe would cause primordial CMBR fluctuations to be erased, whereas the processed γ-ray spectrum could constitute the diffuse γ-ray background at the present epoch. They concluded that this self-consistent model could evade the upper limit on Ω_b given by the observed deuterium abundance and a SBBN scenario since deuterium and ^{3}He would have been produced, at least in part, in the ^{4}He-photodisintegration process. For typical models they produce a fraction $f_{^2\mathrm{H}}^{\mathrm{photo}} \simeq 50\%$ of the total primordial deuterium abundance by ^{4}He photodisintegration. Clearly, this fraction is in conflict with the limit of Eq. (9) and would result in too high $(^3\mathrm{He}/^2\mathrm{H})_p$ ratios [26].

We can also constrain the fraction $f_{(^2\mathrm{H}+^3\mathrm{He})}$ which can be contributed to the total sum of the primordial deuterium and ^{3}He abundances by ^{4}He photodisintegration. This parameter is limited by

$$f_{(^2\mathrm{H}+^3\mathrm{He})}^{\mathrm{photo}} \lesssim 35\text{–}55\,\% \ . \tag{10}$$

Any annihilating topological defects or decaying particles abundant enough to initiate an epoch of cascade nucleosynthesis such that more than 35% of the presently observed abundance sum of $(^2\mathrm{H}+^3\mathrm{He})$ is contributed by this cascade nucleosynthesis are subject to constraint. The limits given in Eq. (10) are a factor 2–3 better than equivalent limits assumed in previous work.

These limits can be put into context by the upper limit on the sum of ^{2}H and ^{3}He inferred from the solar system data and chemical evolution models by Geiss [20]:

$$\left(\frac{^2\mathrm{H} + ^3\mathrm{He}}{\mathrm{H}}\right) \lesssim 1.1 \times 10^{-4} \ . \tag{11}$$

In Fig. 1 we show constraints from ^{4}He photodisintegration on the maximum allowed energy release as a function of redshift. To produce this figure we have used Eq. (11) and the upper end of the range given in Eq. (10). For comparison we show analogous limits from possible distortions of the CMBR background. These are taken from Ref. [12]. It is seen that over a wide range of redshifts the limits from ^{4}He photodisintegration are more stringent than the limits from CMBR distortions. Also shown are constraints from the diffuse γ-ray background which result from the generic cascade spectrum (see Sec. III D).

We can understand the generic features of the limits derived from ^{4}He photodisintegration and shown in Fig. 1 by the solid line from the following considerations. Any radiation injected above the pair-production threshold $E_{\max}$ results effectively instantaneously in a photon spectrum which may be approximated by

$$\frac{dN_\gamma}{dE_\gamma} \approx \begin{cases} \frac{E_0}{E_{\max}^2}\left(\frac{E_\gamma}{E_{\max}}\right)^{-3/2}, & E_\gamma \lesssim E_{\max} \ , \\ 0, \quad E_\gamma \gtrsim E_{\max} \ . \end{cases} \tag{12}$$

Here dN_γ/dE_γ is the number of γ rays per unit energy produced by a total injected energy E_0. The two main processes consuming these second generation γ rays are then the Bethe-Heitler electron-positron pair production on protons and nuclei $(\gamma + p \mapsto p + e^+ + e^-)$ and the

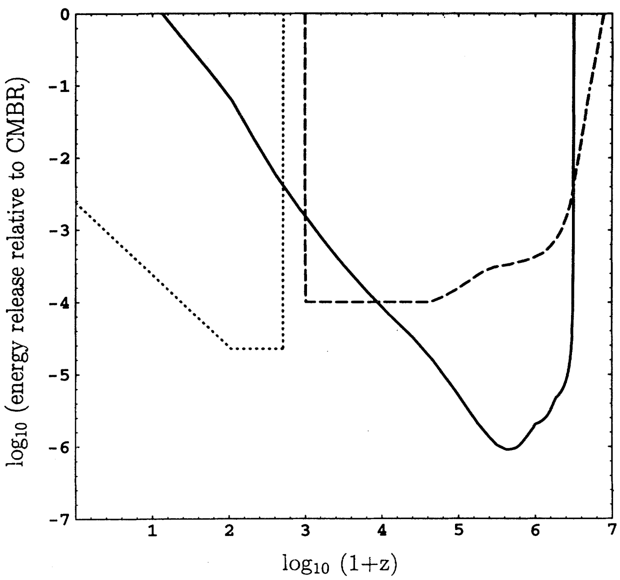

FIG. 1. Maximal energy release in units of the CMBR energy density allowed by the constraints from the observed γ-ray background at 200 MeV (dotted curve), CMBR distortions (dashed curve, from Ref. [13]), and ^{4}He photodisintegration as a function of redshift z. These bounds apply for instantaneous energy release at the specified redshift epoch.

photodisintegration of ^{4}He ($\gamma+{}^4\mathrm{He}\mapsto{}^3\mathrm{He}+n$, etc.). It is a good approximation to only consider the effects of the second generation γ rays since any subsequent generations (produced by $e^\pm + \gamma_{\rm CMBR} \mapsto e^\pm + \gamma$) are generally much softer and may fall below the photodisintegration threshold. The number of ^{3}He nuclei per injected energy produced by the photodisintegration of ^{4}He is then roughly

$$\frac{\Delta N_{^3\mathrm{He}}}{E_0} \approx \sigma_\gamma \Delta t_{\rm BH} n_{^4\mathrm{He}}(z) \int_{E_{\rm th}}^{\infty} \frac{1}{E_0}\frac{dN_\gamma}{dE_\gamma} dE_\gamma \ , \qquad (13)$$

where

$$\Delta t_{\rm BH} \approx [\sigma_{\rm BH} n_{^1\mathrm{H}}(z)]^{-1} \ , \qquad (14)$$

is a typical γ-ray survival time before the γ rays Bethe-Heitler pair produce. In these expressions $n_{^4\mathrm{He}}$ and $n_{^1\mathrm{H}}$ are the ^{4}He- and proton-number densities and $\sigma_\gamma, \sigma_{\rm BH}$ are the cross sections for ^{4}He-photodisintegration and Bethe-Heitler pair production, respectively. Note that we have implicitly assumed that the universe is opaque to γ rays and so our considerations apply only to redshifts $z \gtrsim 10^3$. The resultant ^{3}He to proton ratio at the present epoch is given by

$$\left(\frac{n_{^3\mathrm{He}}}{n_{^1\mathrm{H}}}\right) \approx \frac{1}{n_{^1\mathrm{H}}(z=0)} \frac{\Delta u(z)}{(1+z)^3} \frac{\Delta N_{^3\mathrm{He}}}{E_0} \ , \qquad (15)$$

where $\Delta u(z)$ is the energy density released at epoch with redshift z and the factor $(1+z)^3$ takes into account that light element densities produced during early times are subsequently diluted by the expansion of the universe. This can be translated into a constraint on the injected energy $\Delta u(z)$ relative to the energy density in the CMBR $u_{\rm CMBR}(z)$ at the present epoch by requiring that the light-element abundances not exceed the limits derived from Eqs. (10) and (11). Using Eqs. (12)–(15) we obtain

$$\frac{\Delta u(z)}{u_{\rm CMBR}(z)} \lesssim \frac{1}{(1+z)} \frac{\sigma_{\rm BH}}{\sigma_\gamma} \left(\frac{n_{^3\mathrm{He}}}{n_{^4\mathrm{He}}}\right) \frac{n_{^1\mathrm{H}}(z=0) E_{\rm max}^{1/2} E_{\rm th}^{1/2}}{u_{\rm CMBR}(z=0)} \approx 5\times10^{-7} \left(\frac{1+z}{10^{5.5}}\right)^{-3/2} \ , \qquad (16)$$

where we have used the approximate values $\sigma_{\rm BH} \approx 20$ mb, $\sigma_\gamma \approx 1$ mb, and $E_{\rm th} \approx 5$ MeV. For redshifts $z \gtrsim 10^6$ the threshold for pair production on CMBR photons, $E_{\rm max}$, falls below the photodisintegration threshold, $E_{\rm th}$, and our rough calculation does not apply.

III. ENERGY INJECTION FROM TOPOLOGICAL DEFECTS AND HIGHEST ENERGY COSMIC RAYS

A. History of energy injection in defect models

It is commonly believed that cosmic rays are produced mostly by first order Fermi acceleration (see e.g., [27,28]) at astrophysical shocks in the presence of magnetic fields. The highest energies seem to be reached in relativistic shocks contained in radiogalaxies and active galactic nuclei (see e.g., [29–32]). The recent observation of cosmic rays above 10^{20} eV by the Fly's Eye [33,34] and Akemo Giant Airshower Array (AGASA) [36,37] experiments, and the experiment at Yakutsk [35,38,39] may, however, not be easily explained by this mechanism [40–42]. Therefore, it has been suggested that such superhigh energetic cosmic rays could have a nonacceleration origin [5,8,41,43–46] as, for example, the decay of supermassive elementary "X" particles associated with gand unified theories (GUT's). These particles could be radiated from topological defects (TD's) formed in the early universe during phase transitions caused by spontaneous breaking of symmetries implemented in these GUT's. This is because TD's, like ordinary or superconducting cosmic strings and magnetic monopoles, on which we will focus in this paper, are topologically stable but nevertheless can release part of their energy in form of these X particles due to physical processes like string collapse or monopole annihilation. The X particles with typical GUT scale masses ($\sim 10^{15}$ GeV) decay subsequently into leptons and quarks. The strongly interacting quarks fragment into a jet of hadrons which results in typically of the order of 10^4–10^5 mesons and baryons. It is assumed that these hadrons then give rise to a substantial fraction of the highest energy cosmic ray (HECR) flux, whereas the contribution from the lepton primary is often approximated to be negligible. It also causes a more or less uniform global energy injection whose spectrum is determined by the cascades produced by the interactions of the primary decay products with various background radiation fields. This energy injection is subject to the constraints from ^{4}He photodisintegration discussed in the previous section as well as to constraints from spectral CMBR distortions and the observed γ-ray background.

The X-particle injection rate dn_X/dt as a function of time t or redshift z usually is parametrized as [44]

$$\frac{dn_X}{dt} \propto t^{-4+p} \ . \tag{17}$$

It is important to note that the effective value of p may depend on the epoch. Given that and using standard cosmological relations for $t(z)$ [47] one can describe the X-particle injection history by introducing the dimensionless function

$$g(z) \equiv \frac{1}{(1+z)^3} \frac{(dn_X/dz)(z)}{t_0 (dn_X/dt)(t_0)} \ , \tag{18}$$

where $t_0 = 2H_0^{-1}/3$ is the age of the universe (we assume a flat universe, $\Omega_0 = 1$, throughout this paper).

For example, for annihilating magnetic monopoles it can be shown [46] that $p = 1$ for $t > t_{\rm eq}$ and $p = 1.5$ for $t < t_{\rm eq}$, where $t_{\rm eq}$ is the time of matter-radiation equality.

As a second example, let us look at collapsing cosmic string loops. These may include ordinary as well as superconducting strings. Let us assume that the history of loops consists of two distinct evolutionary stages. We will see below that such a schematic representation can be used for both superconducting strings and ordinary strings. In the first stage the loop slowly radiates gravitational radiation with a power $\sim 100G\mu^2$. Here, G is Newton's constant and $\mu \simeq v^2$ is the energy per unit length of the string in terms of the GUT symmetry breaking scale v. This will decrease the loop length $L(t)$ at an effective rate $v_g \sim 100G\mu$:

$$L(t) = L_b - v_g(t - t_b) \ . \tag{19}$$

In this expression t_b and L_b denote the birth time and the loop length at birth, respectively. Numerical string simulations [48–50] suggest that loops are born with a typical length $L_b = \alpha t_b$ with α being a dimensionless constant which can be as small as a few times v_g. To simplify the calculation we will assume that all loops are born with the same length L_b.

Note that the gravitational radiation associated with this first stage of string loop evolution should not have any effects on CMBR distortions, the diffuse γ-ray background, or result in photodisintegration of ^{4}He. The existence of gravitational radiation during the epoch of primordial nucleosynthesis, however, can effect abundance yields by changing the cosmic expansion rate [51]. For symmetry breaking scales $v \lesssim 10^{16}$ GeV as discussed in this paper this effect is negligible.

Once the loop enters the second evolutionary stage gravitational radiation becomes a subdominant energy loss mechanism. The loop starts to collapse at a rate which grows considerably beyond the gravitational rate v_g by radiating other forms of energy, one of them being X particles. The decay products of these X particles may then contribute to the HECR flux observed at the present epoch. We schematically assume here that during this second evolutionary phase a fraction f of the total energy in loops smaller than a certain critical length scale, $L_c(t)$, is instantaneously released in form of X particles. This is a good approximation as long as the time which loops spend in their second evolutionary phase is short compared to the cosmic time t. By choosing an appropriate value for f this schematic treatment can in an effective way also account for situations where energy release is complicated, for example, by relativistic string motion. Denoting the birth rate of closed string loops per unit volume being chopped off of the string network at birth time t_b by $(dn_b/dt)_{t_b}$ we can then write the rate of X-particle production per unit volume as

$$\frac{dn_X}{dt}(t) = f \left.\frac{dn_b}{dt}\right|_{t_b} \frac{dt_b}{dt} \left[\frac{R(t_b)}{R(t)}\right]^3 \frac{\mu L_c(t)}{m_X} \ . \tag{20}$$

Here $R(t)$ is the cosmic scale factor and $m_X = gv$ is the X-particle mass in terms of the symmetry breaking scale v and the Yukawa coupling $g(g \lesssim 1)$. Furthermore, (dt_b/dt) takes account of the time delay between the birth of a string loop at time t_b and the final phase of X-particle evaporation at later time t. Finally, the factor $[R(t_b)/R(t)]^3$ accounts for dilution due to the cosmic expansion between t_b and t. If the string network exhibits scaling behavior the birth rate of closed string loops can be written as [44]

$$\left.\frac{dn_b}{dt}\right|_{t_b} = \frac{\beta}{t_b^4} \ , \tag{21}$$

where β is a dimensionless constant which is approximately related to α by the relation $\alpha\beta \sim 0.1$ [52].

The possible existence of superconducting cosmic strings within certain GUT's was first proposed by Witten [6]. Ostriker, Thomson, and Witten [7] (OTW) discussed quite severe potential cosmological consequences and also suggested that these objects might contribute to the ultrahigh energy cosmic ray flux. This was further pursued by Hill, Schramm, and Walker [8] who mainly investigated fermionic superconducting string loops which could produce HECR by ejecting superheavy fermion pairs as their length becomes smaller than the (in general time dependent) so called saturation length $L_s(t)$. In addition, these loops radiate electromagnetic waves with a power $\simeq \mu[gL_s(t)/L]^2$. Therefore, as long as $L_s(t)$

is only weakly time dependent, electromagnetic energy loss will eventually dominate the gravitational radiation power $v_g\mu$ for small L. We define the transition time $t_{\rm tr}$ as the cosmic time where these two loss rates are equal for loops born at that time with a length $L_b = \alpha t_{\rm tr}$;

$$\left[\frac{gL_s(t_{\rm tr})}{\alpha t_{\rm tr}}\right]^2 = v_g \ . \qquad (22)$$

With respect to the schematic scenario described above two cosmic epochs have then to be considered for superconducting cosmic strings. For $t \lesssim t_{\rm tr}$, all existing loops are radiating dominantly in electromagnetic and/or X-particle radiation. In this case the loops have not experienced a gravitational radiation dominated energy loss phase, but rather have directly entered the phase of comparatively fast collapse at birth. We can therefore approximate $t_b \simeq t$ and the critical length $L_c(t)$ is the minimum of the birth length, $L_b(t)$, and the saturation length, $L_s(t)$. In contrast, for cosmic times $t > t_{\rm tr}$ the epoch at which a string loop reaches its second evolutionary stage is primarily determined by gravitational energy loss, $t_b \simeq (v_g/\alpha)t$, and $L_c(t) = L_s(t)$. Using Eqs. (20) and (21) this leads to the following time dependence of the X-particle injection rate:

$$\frac{dn_X}{dt} \propto \begin{cases} t^{-4}\left[\frac{R(v_g t/\alpha)}{R(t)}\right]^3 L_s(t) & \text{if } t > t_{\rm tr} \ , \\ t^{-4}{\rm Min}[L_s(t), \alpha t] & \text{if } t < t_{\rm tr} \ . \end{cases} \qquad (23)$$

The saturation length for superconducting strings depends on the intergalactic magnetic field history [8] and is therefore strongly model dependent. OTW originally considered an intergalactic field whose energy density scales like the CMBR energy density. In this scenario the saturation length is roughly constant in time and can be written as

$$L_s(t) \sim {\rm const} \sim 10 \left(\frac{B_0}{10^{-9}\ {\rm G}}\right)\left(\frac{\lambda_0}{1\ {\rm Mpc}}\right)^2\left(\frac{v}{10^{15}\ {\rm GeV}}\right)^{-1/3} g^{-1}\,\alpha^{2/3}\,{\rm pc} \ , \qquad (24)$$

where B_0 and λ_0 are strength and coherence length of the intergalactic field today. Using Eq. (22), the transition time $t_{\rm tr}$ which separates the two string evolution epochs is in terms of redshift $z_{\rm tr}$ given by

$$z_{\rm tr} = 4.78 \times 10^3 \left(\frac{B_0}{10^{-9}\ {\rm G}}\right)^{-1/2}\left(\frac{\lambda_0}{1\ {\rm Mpc}}\right)^{-1}\left(\frac{v}{10^{15}\ {\rm GeV}}\right)^{2/3}\alpha^{1/6} \ . \qquad (25)$$

For the calculations performed in the following we will use $z_{\rm tr} = 2 \times 10^3$. Since we will match the two functional time dependences in Eq. (23) at $t = t_{\rm tr}$ and since ${\rm Min}[L_s(t), \alpha t] = \alpha t$ for $t \ll t_{\rm tr}$ we will use $dn_x/dt \propto t^{-3}$ for all $t < t_{\rm tr}$ for a lower bound on energy injection. Furthermore, we will neglect the time dependence coming from the factor $[R(v_g t/\alpha)/R(t)]^3$ in Eq. (23) in case $t_b = v_g t/\alpha < t_{\rm eq}$ and $t > t_{\rm eq}$. A more detailed treatment would have to take into account the chronological order of t_b, t, and $t_{\rm eq}$ as well as the finite collapse time of a string loop in its second evolutionary stage. This would be model dependent via the parameters from Eqs. (24) and (25). It is, however, easily seen that such effects lead to X-particle injection rates which can only be larger at early times than the injection rates of our simplified treatment Eq. (23). Our calculations will therefore give us conservatively low estimates for the total energy release in X particles. Within these approximations Eq. (23) is of the form of Eq. (17) with $p = 0$ for $t \gtrsim t_{\rm tr}$ and $p = 1$ for $t \lesssim t_{\rm tr}$.

In principle, for superconducting strings the energy radiated in form of X particles is determined by the model. In Ref. [53] it was shown that the ultrahigh energy particles are absorbed in the strong magnetic field produced by the electric current in the string loops. Instead, it was suggested that most of the string energy would be liberated in the form of neutrinos [54]. An absolute flux calculation shows that even without these effects in the OTW scenario, where $L_s(t)$ is approximately constant in time, it is barely possible to produce the observed HECR flux for reasonable model parameters. It has been shown [8] that in scenarios where $L_s(t)$ grows with time the saturation length at the present epoch, $L_s(t_0)$, has necessarily to be smaller than the $L_s(t_0)$ in the OTW scenario. Such scenarios would, for example, be given when intergalactic magnetic fields are increased by dynamo effects. In this case it follows from Eq. (23) that the HECR flux at the present epoch cannot be produced by superconducting cosmic strings even when $f = 1$. In the opposite case [$L_s(t)$ decreases with time] scenarios are conceivable where an $f \lesssim 1$ can reproduce the observed HECR flux. However, for a given universal HECR flux more energy would have been injected outside of the strong magnetic field region in the past compared to the OTW scenario. Therefore, if too much energy tends to be injected within the OTW scenario, as will be shown to be the case below, the other scenarios are also unlikely to be able to explain the observed HECR flux.

In the case of ordinary strings it has been shown that well known physical processes like cusp evaporation are not capable of producing detectable cosmic ray fluxes [10,55]. It has, however, been suggested that a small fraction f of all loops could be formed in states which would lead to their total collapse within one oscillation period after formation [45]. The total energy μL_b in these kinds of loops would be released in form of X particles. Then, $t_b \sim t$ and $L_c(t) = L_b \sim \alpha t$, and Eq. (21) yields

$$\frac{dn_X}{dt} = f\alpha\beta\mu m_X^{-1} t^{-3} , \tag{26}$$

which is of the form of Eq. (17) with $p = 1$. Recently, there has been a claim [56] that loops in high-harmonic states are likely to self-intersect and decay into smaller and smaller loops, finally releasing their energy in relativistic particles. Equation (26) would be a reasonable good approximation also in this case.

Up to now we have only considered the functional form of the X-particle injection rate dn_X/dt up to an absolute normalization. If we *assume* that HECR are produced by decaying X particles radiated from topological defects we can normalize to the differential HECR flux $j_{\rm HECR}(E)$ observed today ($t = t_0$) at a fixed energy $E = E_{\rm obs}$. In these models one expects to observe mainly γ rays at energies $E \gtrsim 10^{20}$ eV [57]. We define the effective X-particle fragmentation function into γ rays, $(dN_\gamma/dx)(x)$ where $x = 2E/m_X$, as the effective differential primary γ-ray multiplicity per injected X particle multiplied by $2/m_X$ [10]. Then the normalization depends on the γ-ray attenuation length $\lambda_\gamma(E)$ and on $(dN_\gamma/dx)(x)$ at $x = 2E_{\rm obs}/m_X$:

$$\begin{aligned}\frac{dn_X}{dt}(t_0) &= \frac{2\pi m_X}{\lambda_\gamma(E_{\rm obs})}\left[\frac{dN_\gamma}{dx}\left(\frac{2E_{\rm obs}}{m_X}\right)\right]^{-1} j_{\rm HECR}(E_{\rm obs}) \\ &\simeq 8.16\times10^{-40}\left(\frac{m_X}{10^{16}\ \mathrm{GeV}}\right)\left(\frac{\lambda_\gamma(E_{\rm obs})}{10\ \mathrm{Mpc}}\right)^{-1}\left(\frac{j_{\rm HECR}(E_{\rm obs})\ \mathrm{GeV\,cm^2\,sec\,sr}}{4\times10^{-31}}\right) \\ &\quad\times\left[\frac{dN_\gamma}{dx}\left(\frac{2E_{\rm obs}}{m_X}\right)\right]^{-1} \mathrm{cm^{-3}\,sec^{-1}} .\end{aligned} \tag{27}$$

In the last expression of Eq. (27) and in the following we have used the numbers for $E_{\rm obs} = 2\times10^{20}$ eV.

Using the parametrization of X-particle injection history, Eq. (18), and the normalization Eq. (27) we are now in a position to derive various constraints on TD models for HECR from limits on energy injection into the universe.

B. Limits from cascade nucleosynthesis

In Ref. [1] the number $N(^3\mathrm{He},\mathrm{D},z)$ of ^{3}He and D nuclei produced via ^{4}He photodisintegration per GeV electromagnetic cascade energy injected into the universe was calculated as a function of redshift z. These functions depend only weakly on h and Ω_b. Therefore, using Eqs. (18) and (27) and assuming that a fraction f_c of the total energy release in high energy particles goes into the cascade one gets

$$\begin{aligned}\left(\frac{^3\mathrm{He}}{\mathrm{H}}\right)_{\rm photo} &\simeq 9.7 f_c\left(\frac{\Omega_b h^2}{0.02}\right)^{-1}\left(\frac{h}{0.75}\right)^{-1}\left(\frac{m_X}{10^{16}\ \mathrm{GeV}}\right)^2\left(\frac{\lambda_\gamma(E_{\rm obs})}{10\ \mathrm{Mpc}}\right)^{-1} \\ &\quad\times\left(\frac{j_{\rm HECR}(E_{\rm obs})\ \mathrm{GeV\,cm^2 sec\,sr}}{4\times10^{-31}}\right)\left[\frac{dN_\gamma}{dx}\left(\frac{2E_{\rm obs}}{m_X}\right)\right]^{-1}\int N(^3\mathrm{He},z)g(z)dz ,\end{aligned} \tag{28}$$

where the integral is performed over the range in Fig. 4 of Ref. [1]. An analogous formula applies for the produced deuterium fraction $(^2\mathrm{H}/\mathrm{H})_{\rm photo}$. Using Eq. (10) and the bound $(^3\mathrm{He}+^2\mathrm{H})/^1\mathrm{H} \le 1.1\times10^{-4}$ we can impose the constraint

$$\left(\frac{^3\mathrm{He}+^2\mathrm{H}}{^1\mathrm{H}}\right)_{\rm photo} \lesssim 5\times10^{-5} . \tag{29}$$

This leads to lower limits on the fragmentation function taken at $x = 2E_{\rm obs}/m_X$ which in the three cases discussed in the previous section read

$$\begin{aligned}\left[\frac{dN_\gamma}{dx}\left(\frac{2E_{\rm obs}}{m_X}\right)\right] &\gtrsim \left\{\begin{array}{ll} 1.4\times10^{5} & \text{for monopole annihilation} \\ 2.0\times10^{6} & \text{for ordinary strings} \\ 1.8\times10^{11} & \text{for the OTW scenario}\end{array}\right\}\times f_c\left(\frac{\Omega_b h^2}{0.02}\right)^{-1}\left(\frac{h}{0.75}\right)^{-1} \\ &\quad\times\left(\frac{m_X}{10^{16}\ \mathrm{GeV}}\right)^2\left(\frac{\lambda_\gamma(E_{\rm obs})}{10\ \mathrm{Mpc}}\right)^{-1}\left(\frac{j_{\rm HECR}(E_{\rm obs})\ \mathrm{GeV\,cm^2\,sec\,sr}}{4\times10^{-31}}\right) .\end{aligned} \tag{30}$$

This has to be compared with expected fragmentation functions in the different defect scenarios. In case of monopoles and ordinary strings this function is mainly determined by the hadronization of the fundamental quarks created in X-particle decays. At HECR energies it is reasonable to assume a power law behavior [45,46]. In superconducting string scenarios the effective spectrum of HECR, which if at all able to leave the high magnetic field region around these strings, could well be altered by interactions with these strong fields. Nevertheless it is still reasonable to assume that at least at HECR energies this spectrum has a power law form.

It can easily be shown that a properly normalized power law fragmentation function $(dN_\gamma/dx)(x) \propto x^{-q}$ ($q > 0$) obeys $(dN_\gamma/dx)(x) \leq 2x^{-2}$ for all $q > 0$. Thus, because of Eq. (30) the OTW scenario is inconsistent with these power law fragmentation functions independent of m_X as long as $f_c \gtrsim 6.9 \times 10^{-3}$. In contrast, the monopole annihilation and ordinary cosmic string scenarios are compatible with reasonable fragmentation functions.

C. Limits from cosmic microwave background distortions

Early nonthermal electromagnetic energy injection can also lead to a distortion of the cosmic microwave background. We focus here on energy injection during the epoch prior to recombination. A comprehensive discussion of this subject was recently given in Ref. [58]. Regarding the character of the resulting spectral CMBR distortions there are basically two periods to distinguish: First, in the range $3 \times 10^6 \simeq z_{\rm th} > z > z_y \simeq 10^5$ between the thermalization redshift $z_{\rm th}$ and the Comptonization redshift z_y, a fractional energy release $\Delta u/u$ leads to a pseudoequilibrium Bose-Einstein spectrum with a chemical potential given by $\mu \simeq 0.71\Delta u/u$. This relation is valid for negligible changes in photon number which is a good approximation for the Klein-Nishina cascades produced by the GUT particle decays we are interested in [58]. Second, in the range $z_y > z > z_{\rm rec} \simeq 10^3$ between z_y and the recombination redshift $z_{\rm rec}$ the resulting spectral distortion is of the Sunyaev-Zel'dovich type [59] with a Compton y parameter given by $4y = \Delta u/u$. The most recent limits on both μ and y were given in Ref. [12]. The resulting bounds on $\Delta u/u$ for instantaneous energy release as a function of injection redshift [13] are shown as the dashed curve in Fig. 1.

Since energy injection by topological defects would be a continuous process it is convenient to define an effective fractional energy release into the CMBR in the following way:

$$\left.\frac{\Delta u}{u}\right|_{\rm eff} \equiv \frac{f_b m_X}{u_0}\int_{z_{\rm rec}}^{z_{\rm th}} \frac{dn_X}{dz}\frac{\xi(z)}{(1+z)^4}dz\ . \tag{31}$$

Here f_b is the fraction of the total energy release in high energy particles which contributes to the CMBR distortion, u_0 is the CMBR energy density today, and $\xi(z)$ is given by 10^{-4} divided by the function shown as the dashed curve in Fig. 1. This effective energy release is constrained to be smaller than 10^{-4} [13]. Similar to Eq. (30) this leads to the lower limits

$$\left[\frac{dN_\gamma}{dx}\left(\frac{2E_{\rm obs}}{m_X}\right)\right] \gtrsim \left\{\begin{array}{ll} 1.2\times 10^5 & \text{for monopole annihilation} \\ 1.5\times 10^5 & \text{for ordinary strings} \\ 1.1\times 10^{10} & \text{for the OTW scenario}\end{array}\right\} \times f_b\left(\frac{h}{0.75}\right)^{-1} \times\left(\frac{m_X}{10^{16}\ \text{GeV}}\right)^2\left(\frac{\lambda_\gamma(E_{\rm obs})}{10\ \text{Mpc}}\right)^{-1}\left(\frac{j_{\rm HECR}(E_{\rm obs})\ \text{GeV cm}^2\,\text{sec sr}}{4\times 10^{-31}}\right)\ . \tag{32}$$

These constraints are less stringent than the constraints Eq. (30) from cascade nucleosynthesis. For the OTW scenario effective power law fragmentation functions are inconsistent with CMBR distortions for $f_b \gtrsim 0.11$.

It should be noted that in the superconducting string scenario there is an additional contribution to the CMBR distortions even if HECR are not produced at all. This contribution comes from the Sunyaev-Zeldovich effect caused by the hot gas produced around the string by emission of electromagnetic radiation before it reaches saturation length and potentially starts to emit HECR. This was discussed in Ref. [7]. Our restriction to distortions caused by HECR alone therefore renders our constraints conservative.

Note that the constraints on the fragmentation function Eqs. (30) and (32) are more stringent for the OTW scenario by a factor of about 10^5 compared to the case of ordinary cosmic strings because for $z > z_{\rm th} \sim 10^3$ the energy release rate from superconducting cosmic strings scales t^{-4} as opposed to t^{-3} for ordinary strings [compare Eqs. (23) and (26)]. In this redshift range both ^{3}He production and CMBR distortions are sensitive to energy injection (see Fig. 1).

D. Limits from the γ-ray background

Electromagnetic cascades which are started at relatively low redshifts z produce an isotropic γ radiation in the observable energy range. Upper limits on the possible flux of ultrahigh energy particles result from a comparison of the predicted and the observed γ-ray flux

[60,62,63].

The limits derived below crucially depend on the assumptions about fragmentation of X particles into the usual particles like protons, pions, photons, electrons, etc., and on the assumption about cosmological evolution of X-particle production [see Eq. (17)]. We shall assume that the fragmentation function for the decay of X particles with mass m_X into particles i ($i = p, \gamma, e$) has the form

$$\frac{2}{m_X}\frac{dN_i}{dx}(x) = \frac{dN_i}{dE_i}(E_i, m_X) = A_i\left(\frac{E_i}{m_X}\right)^{-(q-1)}\frac{1}{E_i}\ , \tag{33}$$

where E_i is the energy of particle i and A_i is a normalization constant. For q we shall focus on the values between $q = 1$ inspired by scaling distribution in inelastic π scattering and $q = 1.32$ according to QCD calculations [5].

As far as evolution is concerned we shall consider two cases: (i) absence of evolution and (ii) the "weak" evolution, as given by Eq. (26) and inspired by the development of a network of cosmic string loops [61].

The strong evolution with $p < 1$ [see Eq. (17)] results in more stringent limits and we shall skip it in this paper.

Let us first turn to the nonevolution case (i). Let the HECR flux observed at $E_{\rm obs} = 2 \times 10^{20}$ eV, $j_{\rm HECR}(E_{\rm obs}) \simeq 4 \times 10^{-31}$ (GeV cm^2 sec sr)$^{-1}$, be caused by protons or γ rays. The generation function for these particles in GeV^{-1} cm^{-3} sec^{-1} can then be found as

$$\Phi_i(E_{\rm obs}) = \frac{4\pi}{\lambda_i(E_{\rm obs})} j_{\rm HECR}(E_{\rm obs})\ , \tag{34}$$

which also leads to Eq. (27). This can be extrapolated to other energies by using the fragmentation function Eq. (33). In Eq. (34) $i = p$ or γ, and $\lambda_i(E_{\rm obs})$ is again the attenuation length for these particles in the CMBR field. From Eqs. (33) and (34) one can then find the total energy production q_i in form of protons, γ rays, and electrons ($i = p, \gamma, e$) in GeV cm^{-3} sec^{-1}.

The energy released in electrons and γ rays (produced directly or through the decay of other particles) goes into electromagnetic cascades (the cascade energy production due to protons is considerably less). Using the usual quark counting one can estimate that about 10% of the total energy release goes into electrons and thus into the cascades. The flux of the cascade protons can then be found as [64]

$$j_\gamma^{\rm cas}(E_\gamma) = \frac{c}{4\pi}\frac{(2/3)H_0^{-1} q_{\rm cas}}{[2 + \ln(E_a/E_x)]E_x^{1/2}} E_\gamma^{-3/2}\ , \tag{35}$$

where E_a and E_x are characteristic cascade energies which for $z = 0$ are given by $E_a \simeq 8 \times 10^4$ GeV and $E_x \simeq 5.1 \times 10^3$ GeV, and $q_{\rm cas}$ is equal to the energy release in the form of electrons and γ rays.

Since observations of the isotropic γ-ray flux in the 100 MeV range yield a power law behavior with an index which is roughly 2 [60,62,63], the most stringent constraints result from a comparison with Eq. (35) at the highest observed energies. Reference [60] reported an upper limit of 7×10^{-8} (MeV cm^2 sec sr)$^{-1}$ at $E_\gamma = 200$ MeV. From Eqs. (34) and (35) we find the cascade flux at this energy to be 8×10^{-8}, 3×10^{-8}, and 9×10^{-9} (MeV cm^2 sec sr)$^{-1}$ for $q = 1.1$, 1.2, and 1.32, respectively, assuming $m_X = 10^{16}$ GeV. For $q = 1.32$ the predicted flux is one order of magnitude less then the observational upper limit.

Let us now go over to the case of evolution (ii). The cascade limit becomes more stringent in this case because the cosmological epochs with large z give no contribution to the presently observed HECR flux at $E \simeq 10^{20}$ eV, while they contribute strongly to the cascade energy density due to the enhanced energy release at earlier times. We shall restrict ourselves to the case of weak evolution here where integration over redshifts results only in a logarithmic factor.

It is easy to understand the existence of a "critical" epoch (with redshift z_c) in our problem. It is defined as $E_\gamma \times (1 + z_c) = E_x(z_c)$, where E_γ is a photon energy at $z = 0$ and $E_x(z_c)$ is the turnover energy of the cascade spectrum at redshift z_c. For $E_\gamma \simeq 200$ MeV one finds $z_c \simeq 100$. If we integrate the evolution function Eq. (26) over the redshift interval between $z = 0$ and $z = z_c$ we obtain

$$j_\gamma^{\rm cas}(E_\gamma) = \frac{c}{4\pi}\frac{H_0^{-1} q_{\rm cas}}{2 + \ln[E_a(z_c)/E_x(z_c)]} \times \frac{\ln(z_c)}{[E_x(0)]^{1/2}} E_\gamma^{-3/2}\ , \tag{36}$$

where $q_{\rm cas}$ is found with the help of Eq. (34) and the fragmentation function Eq. (33) using the energy transfer into p, γ, and e at large redshifts.

For $q = 1.1$, 1.2, and 1.32, the flux Eq. (36) at $E_\gamma \simeq 200$ MeV is numerically 8×10^{-6}, 3×10^{-6}, and 9×10^{-7} (MeV cm^2 sec sr)$^{-1}$, respectively. For X-particle masses different from $m_X = 10^{16}$ GeV these fluxes have to be multiplied by $(m_X/10^{16}\ \text{GeV})^{2-q}$. These numbers are considerably higher than the upper limit 7×10^{-8} (MeV cm^2 sec sr)$^{-1}$ as long as m_X is not much smaller than 10^{16} GeV. These considerations can be translated into the lower limit $q \gtrsim 1.67$ for the index of an assumed power law injection. Recent measurements suggest that the isotropic γ-ray flux might continue to fall roughly as E_γ^{-2} up to 10 GeV [62,63]. In that case the upper limits on $q_{\rm cas}$ might become more stringent by about one order of magnitude and the lower limit on q might come near to 2. In Fig. 2 we compare the γ-ray flux predicted by Eq. (26) for $q = 1.7$ and $m_X = 10^{14}$ GeV with the observational fluxes.

Note that Chi *et al.* [65] derived similar limits by considering cascade development in the CMBR and in the infrared and starlight fields. These limits depend to some extent on the history and intensity of these less well known backgrounds. However, in the case of "weak evolution" of TD's considered here the comparatively strong injection at high redshifts leads to cascading probably mostly in the CMBR, whereas the authors of Ref. [65] were more concerned with low redshift injection where these other backgrounds are more important. Recently Protheroe and Johnson [66] considered limits on energy

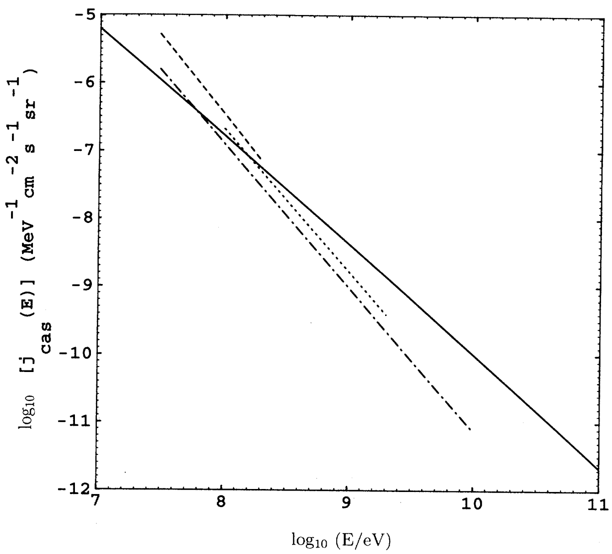

FIG. 2. The generic cascade spectrum (solid line) predicted by the weak evolution scenario for power law injection with an index $q = 1.7$ and $m_X = 10^{14}$ GeV. This scenario serves as an example for the case of HECR production by a network of ordinary cosmic strings. Also shown are observational upper limits from Ref. [60] (short dashed line), Ref. [62] (dash-dotted line), and Ref. [63] (dotted line).

injection from discrete sources and their dependence on the extragalactic magnetic field strength. We belive that these constraints are not directly comparable to our results since we considered uniformly distributed sources in our paper.

As a conclusion we claim that for a fragmentation function of the form of Eq. (33) with reasonable values for $q, 1 \lesssim q \lesssim 1.32$, the explanation of observed HECR at $E \gtrsim 10^{20}$ eV as protons or γ rays from the decay of GUT scale X particles with $m_X \simeq m_{\rm GUT} \simeq 10^{16}$ GeV is incompatible even with the "weak" cosmological evolution of their production. The non-evolution case is not severally constrained by these arguments.

IV. CONCLUSIONS

We have discussed limits on cosmic high energy particle injection derived from ^{4}He photodisintegration, CMBR distortions, and the diffuse γ-ray background. We have found that the nucleosynthesis limits give the most stringent constraints for epochs with redshift $z \gtrsim 5 \times 10^3$ whereas at lower redshifts particle injection is predominantly limited by its contribution to the diffuse γ-ray background (see Fig. 1). These constraints were applied to topological defects potentially radiating supermassive GUT scale ("X") particles which subsequently decay into high energy leptons and hadrons. The history of high energy particle injection is more or less determined within these defect models. The model dependent parameters to be fixed are the number density of X particles radiated within unit time and the effective fragmentation function for the decay products of these X particles. We have assumed that the flux of these decay products contributes significantly to the present day observed HECR flux. This allowed us to formulate our constraints as lower limits on the fractional energy release at HECR energies ($\simeq 10^{20}$ eV) which is mainly determined by the γ-ray fragmentation function. We have found that for reasonable γ-ray fragmentation functions superconducting strings of the type considered in [8] cannot explain the HECR flux without violating at least the bound coming from ^{4}He photodisintegration. In contrast, magnetic monopole and ordinary cosmic string models producing observable HECR fluxes are most severely constrained, but not yet ruled out, by their contribution to the diffuse γ-ray background.

In the second part of the paper we have studied the possibility that the presently observed deuterium has been produced by an epoch of ^{4}He photodisintegration subsequent to a standard nucleosynthesis scenario. Such an epoch may have been initiated by the decay of particles, the annihilation of topological defects, or, in general, the production of energetic γ rays by any source. We have found that only a small fraction ($\lesssim 10\%$) of the observed deuterium may have its origin in the process of ^{4}He photodisintegration since, otherwise, anomalously large primordial (^{3}He/^{2}H) ratios would result. A large fraction of the primordial deuterium contributed by this process would require either standard assumptions of chemical evolution to break down or the existence of γ-ray sources in the early universe which radiate with extremely "soft" γ-ray energy spectra. We have shown that a scenario which employs massive black holes to reprocess the light element abundances from a standard big bang nucleosynthesis process [11] is in conflict with ^{2}H and ^{3}He observations. We have also used the anomaly in the (^{3}He/^{2}H) ratios produced during ^{4}He photodisintegration to slightly tighten constraints on the abundances and parameters of decaying particles and topological defects.

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