



HELIUM PHOTODISINTEGRATION  
AND NUCLEOSYNTHESIS:  
IMPLICATIONS FOR TOPOLOGICAL DEFECTS,  
HIGH ENERGY COSMIC RAYS, AND  
MASSIVE BLACK HOLES

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ABSTRACT

We consider the production of  ${}^3\text{He}$  and  ${}^2\text{H}$  by  ${}^4\text{He}$  photodisintegration initiated by non-thermal energy releases during early cosmic epochs. We find that this process cannot be the predominant source of primordial  ${}^2\text{H}$  since it would result in anomalously high  ${}^3\text{He}/\text{D}$  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  $\gamma$ -rays with energies above  $10^{20}$  eV. The constraints on these models derived from  ${}^4\text{He}$ -photodisintegration are compared to corresponding limits from spectral distortions of the cosmic microwave background radiation (CMBR) and from the observed diffuse  $\gamma$ -ray background. It is shown that for reasonable primary particle injection spectra superconducting cosmic strings, unlike ordinary strings or annihilating monopoles, cannot produce the HECR flux at the present epoch without violating at least the  ${}^4\text{He}$ -photodisintegration bound. The constraint from the diffuse  $\gamma$ -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.

# 1 Introduction

In this paper we consider various constraints inferred from the possible photodisintegration of  ${}^4\text{He}$  in the early universe. Following Protheroe, Stanev, and Berezhinsky [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, 6, 7, 8, 9, 10]. Our focus here will be particularly on constraining the latter scenario. It has also been suggested that  ${}^4\text{He}$ -photodisintegration in the early universe could be a production mechanism for the observed light-element abundances of deuterium and  ${}^3\text{He}$  [11]. In this work we will study the feasibility of such a scenario and show that the  $({}^3\text{He}/{}^2\text{H})$  ratio poses a problem to it. We will show that photodisintegration yields  $({}^3\text{He}/{}^2\text{H}) \gg 1$  and since  ${}^2\text{H}$  is destroyed and  ${}^3\text{He}$  increases with evolution, measures of  $({}^3\text{He}/{}^2\text{H})$  place severe constraints on photodisintegration.

Nonthermal energy releases at high redshifts may leave various observable signatures. The cosmic microwave background radiation (hereafter, 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 re-ionization), 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  $\gamma$ -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  $\gamma$ -rays on protons and  ${}^4\text{He}$  is rare so that the universe becomes transparent to  $\gamma$ -rays with energies below  $E_{\text{max}}$ . Here the energy  $E_{\text{max}}$  is

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

where  $T$  is the CMBR temperature and  $m_e$  is the electron mass.  $E_{\text{max}}$  is related to the threshold energy for  $e^+e^-$ -pair creation by high-energy  $\gamma$ -rays scattering off CMBR-photons. Any radiation with energies above this threshold is effectively instantaneously “recycled” by pair production ( $\gamma\gamma_{\text{CMBR}} \rightarrow e^+e^-$ ) and inverse Compton scattering of the created electrons and positrons ( $e\gamma_{\text{CMBR}} \rightarrow e\gamma$ ). These processes yield a degraded  $\gamma$ -ray spectrum with generic energy dependence  $\propto E_\gamma^{-1.5}$  considerably below  $E_{\text{max}}$  before steepening and finally cutting off at  $E_{\text{max}}$  [3]. Significant energy releases in form of high-energy  $\gamma$ -rays and charged particles at epochs with redshifts below  $z \simeq 300 - 1000$  may therefore produce a present day  $\gamma$ -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  ${}^4\text{He}$  and the concomitant production of deuterium and  ${}^3\text{He}$ . The injection of high-energy particles and  $\gamma$ -rays above the energy threshold  $E_{\text{max}}$  will initiate an epoch of cascade nucleosynthesis subsequent to the epoch of standard primordial nucleosynthesis at  $T \sim 100$  keV. The abundance yields of  ${}^4\text{H}$  and  ${}^3\text{He}$  produced by  ${}^4\text{He}$ -photodisintegration during cascade nucleosynthesis are quite

independent from the primary  $\gamma$ -ray and charged particle energy spectra. Deuterium and  $^3\text{He}$  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 Berezhinsky [1]. The nucleosynthesis limits on the release of energy into the primordial gas can be up to a factor of  $\sim 100$  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  $^4\text{He}$  is inefficient. This is because the energy threshold for pair production falls below the energy threshold for  $^4\text{He}$ -photodisintegration,  $E_{\text{max}} \lesssim E_{th}^{^4\text{He}}$ . The best nucleosynthesis limits on decaying particles and annihilating topological defects in the cosmic temperature range  $1 \text{ keV} \lesssim T \lesssim 10 \text{ keV}$  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  $\gamma\gamma_{\text{CMBR}}$  pair production for  $\gamma$ -energies  $E_\gamma \gtrsim 10^5 \text{ GeV}$  or when there is a significant hadronic decay channel for a massive decaying particle or topological defect. These antinucleons can then annihilate on  $^4\text{He}$  thereby producing approximately equal amounts of  $^2\text{H}$  and  $^3\text{He}$  [17]. We will, however, not further pursue this idea here.

For temperatures above  $T \simeq 1 \text{ keV}$  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  $^4\text{He}$ -mass fractions ( $1 \text{ MeV} \gtrsim T \gtrsim 100 \text{ keV}$ ) and/or increased  $^2\text{H}$  and  $^3\text{He}$ -abundances ( $100 \text{ keV} \gtrsim T \gtrsim 10 \text{ keV}$ ; [18]). It has been suggested that a combination of  $^4\text{He}$ -hadrodestruction and  $^2\text{H}, ^3\text{He}$ -photodestruction induced by a late-decaying particle ( $T \sim 3 \text{ keV}$ ) 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,  $\Omega_b$  [19].

The observational signatures of such scenarios are primordial isotope ratios of  $(^3\text{He}/^2\text{H}) \simeq 2-3$  and  $^6\text{Li}/^7\text{Li} \sim 1$ , 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 life times and hadronic branching ratios, these models would overproduce  $^2\text{H}$  and  $^3\text{He}$  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  $(^3\text{He}/^2\text{H})$  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 Section 2 we briefly review the observationally inferred light-element abundances of  $^2\text{H}$  and  $^3\text{He}$ . We then consider  $^4\text{He}$ -photodisintegration scenarios and their compatibility with the observations. In Section 3 we study the effects of possible energy injection by superconducting strings, ordinary strings, and magnetic monopoles on the primordial  $^2\text{H}$  and  $^3\text{He}$

abundances, the distortions of the CMBR-blackbody, and the diffuse  $\gamma$ -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 Section 4. Throughout this paper we will mostly use  $c = \hbar = 1$ .

## 2 Constraints on ${}^4\text{He}$ -Photodisintegration as the predominant Source of Primordial Deuterium

In this section we investigate scenarios which have  ${}^4\text{He}$ -photodisintegration as an efficient production mechanism of the light-element abundances of deuterium and  ${}^3\text{He}$ . In this study we are naturally led to consider the primordial ratio of  $({}^3\text{He}/{}^2\text{H})_p$ . This is because the ratio of these light isotopes emerging from the big bang nucleosynthesis process,  $({}^3\text{He}/{}^2\text{H})_{\text{BBN}}$ , is quite different from that emerging from the  ${}^4\text{He}$ -photodisintegration,  $({}^3\text{He}/{}^2\text{H})_{\text{photo}}$ . In particular, we expect generic isotope ratios of  $({}^3\text{He}/{}^2\text{H})_{\text{BBN}} \lesssim 1$ , and  $({}^3\text{He}/{}^2\text{H})_{\text{photo}} \gg 1$ . We will show that this fact can be used to severely constrain the photodisintegration of  ${}^4\text{He}$  as the principal source of primordial deuterium. We will also show that the observationally inferred abundances of  ${}^2\text{H}$  and  ${}^3\text{He}$  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  $({}^3\text{He}/{}^2\text{H})$ -ratio is thought to come from solar system observations of  ${}^3\text{He}$  abundances. Geiss [20] reanalyzed the existing data and inferred for the abundances of deuterium and  ${}^3\text{He}$  at the time of solar system formation

$$\begin{aligned}
 1.2 \times 10^{-5} &\lesssim \left(\frac{{}^3\text{He}}{\text{H}}\right)_{\oplus} \lesssim 1.8 \times 10^{-5}, \\
 1.6 \times 10^{-5} &\lesssim \left(\frac{{}^2\text{H}}{\text{H}}\right)_{\oplus} \lesssim 3.3 \times 10^{-5}, \\
 0.34 &\lesssim \left(\frac{{}^3\text{He}}{{}^2\text{H}}\right)_{\oplus} \lesssim 1.13.
 \end{aligned} \tag{2}$$

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

Deuterium is the most fragile of the light isotopes. It is easily destroyed during the pre-main sequence evolutionary stage of stars via  ${}^2\text{H}(p,\gamma){}^3\text{He}$ . 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  ${}^3\text{He}$  is less clear. It is known that  ${}^3\text{He}$  is destroyed to some extent in massive stars ( $M \gtrsim 5 - 8M_{\odot}$ ), whereas low-mass stars ( $M \lesssim 1 - 2M_{\odot}$ ) may be net producers of  ${}^3\text{He}$ .

This theory is supported by the observations of  ${}^3\text{He}$  abundances in planetary nebulae. It is certainly very reasonable to assume that standard chemical evolution models can only increase the primordial  $({}^3\text{He}/{}^2\text{H})_p$ -ratio.

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

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

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

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

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

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

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

for a wide range of fractional contributions of baryons to the closure density,  $\Omega_b$ , Hubble parameters  $H_0$  in units of  $100 \text{ km sec}^{-1} \text{ Mpc}^{-1}$ ,  $h$ , and epochs of energy injection. This is because in  ${}^4\text{He}$ -photodisintegration the effective cross sections for the two-nucleon photoabsorption processes [ ${}^4\text{He}(\gamma, \text{pn}){}^2\text{H}$  and  ${}^4\text{He}(\gamma, {}^2\text{H}){}^2\text{H}$ ] are roughly ten times smaller than the effective cross sections for the single-nucleon photoabsorption processes [ ${}^4\text{He}(\gamma, \text{p}){}^3\text{H}$  and  ${}^4\text{He}(\gamma, \text{n}){}^3\text{He}$ ] [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  $\gamma\gamma_{\text{CMBR}} \rightarrow e^-e^+$  pair creation. Pair creation and inverse Compton scattering will then yield a generic  $\gamma$ -ray spectrum with energy dependence  $\propto E_\gamma^{-1.5}$  below  $E_{\text{max}}/2$  and  $\propto E_\gamma^{-5}$  above before cutting off at  $E_{\text{max}}$ . These  $\gamma$ -rays can be effective in photodisintegrating  ${}^4\text{He}$  where the competing process is the consumption of  $\gamma$ -rays by Bethe-Heitler pair production on hydrogen and helium.

When radiation is injected below  $E_{\text{max}}$  the  $\gamma$ -rays may have a spectrum quite different from the behavior  $\propto E_\gamma^{-1.5}$  depending on the actual  $\gamma$ -ray source. In principle, it is then

conceivable to photodisintegrate  ${}^4\text{He}$  in such a way that isotope ratios of  $({}^3\text{He}/{}^2\text{H}) \simeq 1$  result. This could be accomplished by a  $\gamma$ -ray source which preferentially radiates above energies of  $E \simeq 100$  MeV but below  $E_{\text{max}}$ . This is because only in the energy range between the  ${}^4\text{He}$ -photodisintegration threshold  $E_{th}^{4\text{He}} = 19.8$  MeV and  $E \simeq 100$  MeV the effective cross section for  ${}^3\text{He}$  production in  ${}^4\text{He}$ -photodisintegration is roughly ten times larger than the effective cross section for  ${}^2\text{H}$  production in this process. For  $\gamma$ -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  $\gamma$ -ray spectrum produced by Bethe-Heitler pair production and inverse Compton scattering. In this case, however, significant deuterium production would require  $\gamma$ -ray fluxes which would exceed the present day diffuse  $\gamma$ -ray background.

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

We can derive limits on the allowed contributions of  ${}^4\text{He}$ -photodisintegration to the primordial  ${}^2\text{H}$  and  ${}^3\text{He}$  abundances. This can be done by employing the solar system  $({}^3\text{He}/{}^2\text{H})$ -isotope ratio from Eq. (2) and assuming that this ratio represents a conservative upper limit on the primordial  $({}^3\text{He}/{}^2\text{H})$ -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  $f_{2\text{H}}^{\text{photo}}$  contributed to the primordial deuterium abundance by  ${}^4\text{He}$ -photodisintegration. A simple calculation of the abundance average then yields

$$f_{2\text{H}}^{\text{photo}} \lesssim \frac{(\frac{{}^3\text{He}}{{}^2\text{H}})_{\odot} - (\frac{{}^3\text{He}}{{}^2\text{H}})_{\text{BBN}}}{(\frac{{}^3\text{He}}{{}^2\text{H}})_{\text{photo}} - (\frac{{}^3\text{He}}{{}^2\text{H}})_{\text{BBN}}} . \quad (6)$$

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

$$f_{2\text{H}}^{\text{photo}} \lesssim 13\% . \quad (7)$$

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  ${}^3\text{He}$ . *This also implies that generic  ${}^4\text{He}$ -photodisintegration scenarios can not be the predominant production mechanism of the primordial  ${}^2\text{H}$  and  ${}^3\text{He}$  light-element abundances.* The stringent limit of Eq. (7) can only be evaded when either there existed an extremely “soft”  $\gamma$ -ray source in the early universe or when generic features of the galactic destruction/production of  ${}^3\text{He}$  and  ${}^2\text{H}$  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 accret material which emits a quasar-like X-ray and  $\gamma$ -ray spectrum they may induce the photodisintegration of  $^4\text{He}$  and the reionization of the universe. The reionization of the universe would cause primordial CMBR fluctuations to be erased, whereas the processed  $\gamma$ -ray spectrum could constitute the diffuse  $\gamma$ -ray background at the present epoch. They concluded that this selfconsistent model could evade the upper limit on  $\Omega_b$  given by the observed deuterium abundance and a SBBN scenario since deuterium and  $^3\text{He}$  would have been produced, at least in part, in the  $^4\text{He}$ -photodisintegration process. For typical models they produce a fraction  $f_{^2\text{H}}^{\text{photo}} \simeq 50\%$  of the total primordial deuterium abundance by  $^4\text{He}$ -photodisintegration. Clearly, this fraction is in conflict with the limit of Eq. (7) and would result in too high  $(^3\text{He}/^2\text{H})_p$  ratios [26].

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

$$f_{(^2\text{H}+^3\text{He})}^{\text{photo}} \lesssim 35\% - 55\%. \quad (8)$$

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  $(^2\text{H}+^3\text{He})$  is contributed by this cascade nucleosynthesis are subject to constraint. The limits given in Eq. (8) 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  $^2\text{H}$  and  $^3\text{He}$  inferred from the solar system data and chemical evolution models by Geiss [20]

$$\left(\frac{^2\text{H} + ^3\text{He}}{\text{H}}\right) \lesssim 1.1 \times 10^{-4}. \quad (9)$$

In Figure 1 we show constraints from  $^4\text{He}$ -photodisintegration on the maximum allowed energy release as a function of redshift. To produce this figure we have used Eq. (9) and the upper range given in Eq. (8). For comparison we show analogous limits from possible distortions of the CMBR background. These are taken from reference [12]. It is seen that over a wide range of redshifts the limits from  $^4\text{He}$  photodisintegration are more stringent than the limits from CMBR distortions. Also shown are constraints from the diffuse  $\gamma$ -ray background which result from the generic cascade spectrum (see section 3.4).

### 3 Energy Injection from Topological Defects and Highest Energy Cosmic Rays

#### 3.1 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, 30, 31, 32]). The recent observation of cosmic rays above  $10^{20}$  eV by the Fly's Eye [33, 34] and AGASA [36, 37] experiments, and the experiment at Yakutsk [38, 39] may, however, not be easily explained by this mechanism [40, 41, 42]. Therefore, it has been suggested that such superhigh energetic cosmic rays could have a non-acceleration origin [5, 8, 41, 43, 44, 45, 46] as, for example, the decay of supermassive elementary "X" particles associated with Grand Unified Theories (GUTs). These particles could be radiated from topological defects (TDs) formed in the early universe during phase transitions caused by spontaneous breaking of symmetries implemented in these GUTs. This is because TDs, 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 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  ${}^4\text{He}$ -photodisintegration discussed in the previous section as well as to constraints from spectral CMBR distortions and the observed  $\gamma$ -ray background.

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

$$\frac{dn_X}{dt} \propto t^{-4+p}. \quad (10)$$

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

$$f(z) \equiv \frac{(dn_X/dz)(z)}{t_0(dn_X/dt)(t_0)}, \quad (11)$$

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  $t > t_{\text{eq}}$  and  $p = 1.5$  for  $t < t_{\text{eq}}$ , where  $t_{\text{eq}}$  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  $v_g \sim 100G\mu$ .

$$L(t) = L_b - v_g(t - t_b). \quad (12)$$

In this expression  $t_b$  and  $L_b$  denote the birth time and the loop length at birth, respectively. Numerical string simulations [48, 49, 50] suggest that loops are born with a typical length  $L_b = \alpha t_b$  with  $\alpha$  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  $\gamma$ -ray background, or result in photodissociation of  ${}^4\text{He}$ . 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  $t$ . 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

$$\frac{dn_X}{dt}(t) = f \left. \frac{dn_b}{dt} \right|_{t_b} \frac{dt_b}{dt} \left[ \frac{R(t_b)}{R(t)} \right]^3 \frac{\mu L_c(t)}{m_X}. \quad (13)$$

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]

$$\left. \frac{dn_b}{dt} \right|_{t_b} = \frac{\beta}{t_b^4}, \quad (14)$$

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

The possible existence of superconducting cosmic strings within certain GUTs was first proposed by Witten [6]. Ostriker, Thomson and Witten [7] (hereafter, 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 HE CR by ejecting superheavy fermion pairs towards the end of their evolution. With respect to the schematic scenario described above two cosmic epochs have then to be considered for superconducting cosmic strings of this type. For cosmic time  $t \lesssim t_{\text{tr}}$ , all existing loops are radiating dominantly in electromagnetic and/or X-particle radiation. In terms of the (in general time dependent) saturation length  $L_s(t)$  the relevant condition which compares electromagnetic and gravitational energy loss rates reads [8]

$$\left[ \frac{gL_s(t)}{\alpha t} \right]^2 > v_g. \quad (15)$$

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  $t > t_{\text{tr}}$  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. (13) and (14) this leads to the following time dependence of the X-particle injection rate:

$$\frac{dn_X}{dt} \propto \begin{cases} t^{-4} \left[ \frac{R(v_g t/\alpha)}{R(t)} \right]^3 L_s(t) & \text{if } t > t_{\text{tr}} \\ t^{-4} \text{Min}(L_s(t), \alpha t) & \text{if } t < t_{\text{tr}} \end{cases} \quad (16)$$

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

$$L_s(t) \sim \text{const.} \sim 10 \left( \frac{B_0}{10^{-9} \text{G}} \right) \left( \frac{\lambda_0}{1 \text{Mpc}} \right)^2 \left( \frac{v}{10^{15} \text{GeV}} \right)^{-1/3} g^{-1} \alpha^{2/3} \text{pc}. \quad (17)$$

where  $B_0$  and  $\lambda_0$  are strength and coherence length of the intergalactic field today. Using Eq (15), the transition time  $t_{\text{tr}}$  which separate the two string evolution epochs is in terms of redshift  $z_{\text{tr}}$  given by

$$z_{\text{tr}} = 4.78 \times 10^3 \left( \frac{B_0}{10^{-9} \text{G}} \right)^{-1/2} \left( \frac{\lambda_0}{1 \text{Mpc}} \right)^{-1} \left( \frac{v}{10^{15} \text{GeV}} \right)^{2/3} \alpha^{1/6}. \quad (18)$$

For the calculations performed in the following we will use  $z_{\text{tr}} = 2 \times 10^3$ . Since we will match the two functional time dependences in Eq. (16) at  $t = t_{\text{tr}}$  and since  $\text{Min}(L_s(t), \alpha t) = \alpha t$  for  $t \ll t_{\text{tr}}$  we will use  $dn_x/dt \propto t^{-3}$  for all  $t < t_{\text{tr}}$  for a lower bound on energy injection. Furthermore, we will neglect the time dependence coming from the factor  $[R(v_g t/\alpha) \cdot R(t)]^3$  in Eq. (16) in case  $t_b = v_g t/\alpha < t_{\text{eq}}$  and  $t > t_{\text{eq}}$ . A more detailed treatment would have to

take into account the chronological order of  $t_b$ ,  $t$  and  $t_{\text{eq}}$  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. (17) and (18). 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. (16). Our calculations will therefore give us conservatively low estimates for the total energy release in X-particles. Within these approximations Eq. (16) is of the form of Eq. (10) with  $p = 0$  for  $t \gtrsim t_{\text{tr}}$  and  $p = 1$  for  $t \lesssim t_{\text{tr}}$ .

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]. We shall demonstrate here 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.(16) that the HECR flux at the present epoch can not be produced by superconducting cosmic strings even when  $f = 1$ . In the opposite case ( $L_s(t)$  growing in 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 in these kinds of loops would be released in form of X-particles. Then,  $t_b \sim t$  and  $L_c(t) = L_b \sim at$  so that

$$\frac{dn_X}{dt} = f\alpha\beta\mu m_X^{-1}t^{-3}, \quad (19)$$

which is of the form of Eq. (10) 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. Eq. (19) 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  $j_{\text{HECR}}(E)$  observed today ( $t = t_0$ ) at a fixed energy  $E = E_{\text{obs}}$ . In these models one expects to observe mainly  $\gamma$ -rays at energies  $E \gtrsim 10^{20}$  eV [57]. We define the effective X-

particle fragmentation function into  $\gamma$ -rays,  $(dN_\gamma/dx)(x)$  where  $x = 2E/m_X$ , as the effective differential primary  $\gamma$ -ray multiplicity per injected X-particle multiplied by  $2/m_X$  [10]. Then the normalization depends on the  $\gamma$ -ray attenuation length  $\lambda_\gamma(E)$  and on  $(dN_\gamma/dx)(x)$  at  $x = 2E_{\text{obs}}/m_X$ :

$$\begin{aligned} \frac{dn_X}{dt}(t_0) &= \frac{2\pi m_X}{\lambda_\gamma(E_{\text{obs}})} \left[ \frac{dN_\gamma}{dx} \left( \frac{2E_{\text{obs}}}{m_X} \right) \right]^{-1} j_{\text{HECR}}(E_{\text{obs}}) \\ &\simeq 8.16 \times 10^{-40} \left( \frac{m_X}{10^{16} \text{ GeV}} \right) \left( \frac{\lambda_\gamma(E_{\text{obs}})}{10 \text{ Mpc}} \right)^{-1} \left( \frac{j_{\text{HECR}}(E_{\text{obs}}) \cdot \text{GeV cm}^2 \text{ sec sr}}{4 \times 10^{-31}} \right) \\ &\times \left[ \frac{dN_\gamma}{dx} \left( \frac{2E_{\text{obs}}}{m_X} \right) \right]^{-1} \text{cm}^{-3} \text{sec}^{-1}. \end{aligned} \quad (20)$$

In the last expression of Eq. (20) and in the following we have used the numbers for  $E_{\text{obs}} = 2 \times 10^{20} \text{ eV}$ .

Using the parametrization of X-particle injection history, Eq. (11), and the normalization Eq. (20) we are now in a position to derive various constraints on TD models for HE CR from limits on energy injection into the universe.

### 3.2 Limits from Cascade Nucleosynthesis

In Ref. [1] the number  $N(^3\text{He}, \text{D}, z)$  of  $^3\text{He}$  and D nuclei produced via  $^4\text{He}$ -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  $\Omega_b$ . Therefore, using Eqs. (11) and (20) and assuming that a fraction  $f_c$  of the total energy release in high energy particles goes into the cascade one gets

$$\begin{aligned} \left( \frac{^3\text{He}}{\text{H}} \right)_{\text{photo}} &\simeq 9.7 f_c \left( \frac{\Omega_b h^2}{0.02} \right)^{-1} \left( \frac{h}{0.75} \right)^{-1} \left( \frac{m_X}{10^{16} \text{ GeV}} \right)^2 \left( \frac{\lambda_\gamma(E_{\text{obs}})}{10 \text{ Mpc}} \right)^{-1} \\ &\times \left( \frac{j_{\text{HECR}}(E_{\text{obs}}) \cdot \text{GeV cm}^2 \text{ sec sr}}{4 \times 10^{-31}} \right) \left[ \frac{dN_\gamma}{dx} \left( \frac{2E_{\text{obs}}}{m_X} \right) \right]^{-1} \int N(^3\text{He}, z) \frac{f(z)}{(1+z)^3} dz, \end{aligned} \quad (21)$$

where the integral is performed over the range in Fig. 4 of Ref. [1]. An analogous formula applies for the produced deuterium fraction  $(^2\text{H}/\text{H})_{\text{photo}}$ . Using Eq. (8) and the bound  $(^3\text{He} + ^2\text{H})/\text{H} \leq 1.1 \times 10^{-4}$  we can impose the constraint

$$\left( \frac{^3\text{He} + \text{D}}{\text{H}} \right)_{\text{photo}} \lesssim 5 \times 10^{-5}. \quad (22)$$

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

$$\left[ \frac{dN_\gamma}{dx} \left( \frac{2E_{\text{obs}}}{m_X} \right) \right] \gtrsim \left\{ \begin{array}{ll} 1.4 \times 10^5 & \text{for monopole annihilation} \\ 2.0 \times 10^6 & \text{for ordinary strings} \\ 1.8 \times 10^{11} & \text{for the OTW scenario} \end{array} \right\} \times f_c \left( \frac{\Omega_b h^2}{0.02} \right)^{-1} \left( \frac{h}{0.75} \right)^{-1}$$

$$\times \left( \frac{m_X}{10^{16} \text{ GeV}} \right)^2 \left( \frac{\lambda_\gamma(E_{\text{obs}})}{10 \text{ Mpc}} \right)^{-1} \left( \frac{j_{\text{HECR}}(E_{\text{obs}}) \cdot \text{GeV cm}^2 \text{ sec sr}}{4 \times 10^{-31}} \right). \quad (23)$$

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 HECD energies it is reasonable to assume a power law behavior [45, 46]. In superconducting string scenarios the effective spectrum of HECD, 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 HECD 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. (23) 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.

### 3.3 Limits from Cosmic Microwave Background Distortions

Early non-thermal 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  $3 \times 10^6 \simeq z_{\text{th}} > z > z_y \simeq 10^5$  between the thermalization redshift  $z_{\text{th}}$  and the Comptonization redshift  $z_y$ , a fractional energy release  $\Delta u/u$  leads to a pseudo-equilibrium 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  $z_y > z > z_{\text{rec}} \simeq 10^3$  between  $z_y$  and the recombination redshift  $z_{\text{rec}}$  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  $\mu$  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:

$$\frac{\Delta u}{u} \Big|_{\text{eff}} \equiv \frac{f_b m_X}{u_0} \int_{z_{\text{rec}}}^{z_{\text{th}}} \frac{dn_X}{dz} \frac{\xi(z)}{(1+z)^4} dz. \quad (24)$$

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. (23) this leads to the lower limits

$$\left[ \frac{dN_\gamma}{dx} \left( \frac{2E_{\text{obs}}}{m_X} \right) \right] \gtrsim \left\{ \begin{array}{ll} 1.2 \times 10^5 & \text{for monopole annihilation} \\ 1.5 \times 10^5 & \text{for ordinary strings} \\ 1.1 \times 10^{10} & \text{for the OTW scenario} \end{array} \right\} \times f_b \left( \frac{h}{0.75} \right)^{-1} \quad (25)$$

$$\times \left( \frac{m_X}{10^{16} \text{ GeV}} \right)^2 \left( \frac{\lambda_\gamma(E_{\text{obs}})}{10 \text{ Mpc}} \right)^{-1} \left( \frac{j_{\text{HECR}}(E_{\text{obs}}) \cdot \text{GeV cm}^2 \text{ sec sr}}{4 \times 10^{-31}} \right)^{-1}.$$

These constraints are less stringent than the constraints Eq. (23) 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.

### 3.4 Limits from the $\gamma$ -ray Background

Electromagnetic cascades which are started at relatively low redshifts  $z$  produce an isotropic  $\gamma$ -radiation in the observable energy range. The flux in this radiation puts an upper limit on the possible flux of ultrahigh energy particles. The most stringent constraint comes from the upper limit to the observed isotropic flux at  $E_\gamma \simeq 200 \text{ MeV}$ , which was reported to be  $7 \times 10^{-8} (\text{MeV cm}^2 \text{ sec sr})^{-1}$  [61].

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. (10)].

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

$$\frac{2}{m_X} \frac{dN_i}{dx}(x) = \frac{dN_i}{dE_i}(E_i, m_X) = A_i \left( \frac{E_i}{m_X} \right)^{-(q-1)} \frac{1}{E_i}. \quad (26)$$

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 pp-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. (19) and inspired by the development of a network of cosmic string loops [62]. The strong evolution with  $p < 1$  [see Eq. (10)] results in more stringent limits and we shall skip it in this paper.

Let us first turn to the non-evolution case (i). Let the HECR flux observed at  $E_{\text{obs}} = 2 \times 10^{20} \text{ eV}$ ,  $j_{\text{HECR}}(E_{\text{obs}}) \simeq 4 \times 10^{-31} (\text{GeV cm}^2 \text{ sec sr})^{-1}$ , be caused by protons or  $\gamma$ -rays. The

generation function for these particles in  $\text{GeV}^{-1} \text{cm}^{-3} \text{sec}^{-1}$  can then be found as

$$\Phi_i(E_{\text{obs}}) = \frac{4\pi}{\lambda_i(E_{\text{obs}})} j_{\text{HECR}}(E_{\text{obs}}). \quad (27)$$

which also leads to Eq. (20). This can be extrapolated to other energies by using the fragmentation function Eq. (26). In Eq. (27)  $i = p$  or  $\gamma$ , and  $\lambda_i(E_{\text{obs}})$  is again the attenuation length for these particles in the CMBR field. From Eqs. (26) and (27) one can then find the total energy production  $q_i$  in form of protons,  $\gamma$ -rays and electrons ( $i = p, \gamma, e$ ) in  $\text{GeV cm}^{-3} \text{sec}^{-1}$ .

The energy released in electrons and  $\gamma$ -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 photons can then be found as [63]

$$j_{\gamma}^{\text{cas}}(E_{\text{obs}}) = \frac{c}{4\pi} \frac{(2/3)H_0^{-1}q_{\text{cas}}}{[2 + \ln(E_a/E_x)]E_x^{1/2}} E_{\gamma}^{-3/2}. \quad (28)$$

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

From Eqs. (27) and (28) we find the cascade flux at  $E_{\gamma} \simeq 200 \text{ MeV}$  to be  $8 \times 10^{-8}$ ,  $3 \times 10^{-8}$ , and  $9 \times 10^{-9} (\text{MeV cm}^2 \text{sec sr})^{-1}$  for  $q = 1.1, 1.2$  and  $1.32$ , respectively, assuming  $m_X = 10^{16} \text{ GeV}$ . These numbers should be compared with the observational upper limit  $7 \times 10^{-8} (\text{MeV cm}^2 \text{sec sr})^{-1}$  [61]. For  $q = 1.32$  the predicted flux is one order of magnitude less.

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 HE CR flux at  $E \simeq 10^{20} \text{ eV}$ , 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  $E_{\gamma}$  is a photon energy at  $z = 0$  and  $E_x(z_c)$  is the turn-over energy of the cascade spectrum at redshift  $z_c$ . For  $E_{\gamma} \simeq 200 \text{ MeV}$  one finds  $z_c \simeq 100$ . If we integrate the evolution function Eq. (19) over the redshift interval between  $z = 0$  and  $z = z_c$  we obtain

$$j_{\gamma}^{\text{cas}}(E_{\text{obs}}) = \frac{c}{4\pi} \frac{H_0^{-1}q_{\text{cas}}}{2 + \ln[E_a(z_c)/E_x(z_c)]} \frac{\ln(z_c)}{[E_x(0)]^{1/2}} E_{\gamma}^{-3/2}, \quad (29)$$

where  $q_{\text{cas}}$  is found with the help of Eq. (27) and the fragmentation function Eq. (26) using the energy transfer into  $p, \gamma$  and  $e$  at large redshifts.

For  $q = 1.1, 1.2$ , and  $1.32$ , the flux Eq. (29) at  $E_{\gamma} \simeq 200 \text{ MeV}$  is numerically  $8 \times 10^{-6}$ ,  $3 \times 10^{-6}$ , and  $9 \times 10^{-7} (\text{MeV cm}^2 \text{sec sr})^{-1}$ , 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  $7 \times 10^{-8} (\text{MeV cm}^2 \text{sec sr})^{-1}$  as long as  $m_X$  is not much smaller than  $10^{16}$  GeV. These considerations can be translated into the lower limit  $q \gtrsim 1.6$  for the index of an assumed power law injection.

Note that Chi *et al.* [64] 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 TDs considered here the comparatively strong injection at high redshifts leads to cascading probably mostly in the CMBR, whereas the authors of Ref. [64] were more concerned with low redshift injection where these other backgrounds are more important.

As a conclusion we claim that for a fragmentation function of the form of Eq. (26) 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  $\gamma$ -rays from the decay of GUT scale X-particles with  $m_X \simeq m_{\text{GUT}} \simeq 10^{16}$  GeV is incompatible even with the “weak” cosmological evolution of their production. The non-evolution case is not severely constrained by these arguments.

## 4 Conclusions

We have discussed limits on cosmic high energy particle injection derived from  $^4\text{He}$  photodisintegration, CMBR distortions and the diffuse  $\gamma$ -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  $\gamma$ -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  $\gamma$ -ray fragmentation function. We have found that for reasonable  $\gamma$ -ray fragmentation functions superconducting strings can not explain the HECR flux without violating at least the bound coming from  $^4\text{He}$ -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  $\gamma$ -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  $^4\text{He}$ -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  $\gamma$ -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  $^4\text{He}$ -photodisintegration since, otherwise, —

anomalously large primordial ( ${}^3\text{He}/{}^2\text{H}$ )-ratios would result. A larger fraction of the primordial deuterium contributed by this process would require either standard assumptions of chemical evolution to break down or the existence of  $\gamma$ -ray sources in the early universe which radiate with extremely "soft"  $\gamma$ -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  ${}^2\text{H}$  and  ${}^3\text{He}$  observations. We have also used the anomaly in the ( ${}^3\text{He}/{}^2\text{H}$ )-ratios produced during  ${}^4\text{He}$ -photodisintegration to slightly tighten constraints on the abundances and parameters of decaying particles and topological defects.

## Acknowledgments

This work was supported by the DoE, NSF and NASA at the University of Chicago, by the DoE and by NASA through grant NAGS-2788 at Fermilab, and by the Alexander-von-Humboldt Foundation. This work was also performed under the auspices of the U.S. Department of Energy by the Lawrence Livermore National Laboratory under contract number W-7405-ENG-48 and DoE Nuclear Theory Grant SF-ENG-48.

## References

- [1] R. J. Protheroe, T. Stanev, and V. S. Berezinsky, Preprint LNGS-94/106 (1994).
- [2] D. Lindley, *MNRAS* **188**, 15p (1979); V. S. Berezinsky, *Nucl. Phys.* **B380**, 478 (1992).
- [3] J. Ellis, G. B. Gelmini, J. L. Lopez, D. V. Nanopoulos, and S. Sarkar, *Nucl. Phys.* **B373**, 399 (1992); see also references therein.
- [4] S. Miyama and K. Sato, *Prog. Theor. Phys.* **59**, 1012 (1978); D. Lindley, *MNRAS* **193**, 593 (1980); T. Rothman and R. Matzner, *Astrophys. Space Sci.* **75**, 229 (1981).
- [5] C. T. Hill, *Nucl. Phys.* **B224**, 469 (1983).
- [6] E. Witten, *Nucl. Phys.* **B249**, 557 (1985).
- [7] J. P. Ostriker, C. Thompson, and E. Witten, *Phys. Lett.* **180**, 231 (1986).
- [8] C. T. Hill, D. N. Schramm, and T. P. Walker, *Phys. Rev.* **D36**, 1007 (1987).
- [9] R. H. Brandenberger, *Nucl. Phys.* **B293**, 812 (1987).
- [10] P. Bhattacharjee, *Phys. Rev.* **D40**, 3968 (1989).
- [11] N. Yu. Gnedin and J. P. Ostriker, *Astrophys. J.* **400**, 1 (1992).
- [12] J. C. Mather *et al.*, *Astrophys. J.* **420**, 439 (1994).

- [13] E. L. Wright *et al.*, *Astrophys. J.* **420**, 450 (1994).
- [14] P. J. E. Peebles, *Physical Cosmology* (Princeton: Princeton Univ. Press, 1971); R. A. Sunyaev and Ya. B. Zel'dovich, *ARA&A* **18**, 537 (1980); A. P. Lightman, *Astrophys. J.* **244**, 392 (1981).
- [15] J. I. Trombka and C. E. Fichtel, *Phys. Rep.* **97**, 173 (1983).
- [16] S. Dimopoulos, R. Esmailzadeh, L. J. Hall, and G. D. Starkman, *Nucl. Phys.* **B311**, 699 (1989).
- [17] F. Balestra *et al.*, *Il Nuovo Cimento* **100A** N.3, 323 (1988); F. Balestra *et al.*, *Phys. Lett.* **B230**, 36 (1989).
- [18] M. H. Reno and D. Seckel, *Phys. Rev.* **D37**, 3441 (1988).
- [19] S. Dimopoulos, R. Esmailzadeh, L. J. Hall, and G. D. Starkman, *Astrophys. J.* **330**, 545 (1988).
- [20] J. Geiss, in *Origin and Evolution of the Elements* ed. N. Prantzos, E. Vangioni-Flam, and M. Casse (Cambridge: Cambridge Univ. Press), 89 (1993).
- [21] K. A. Olive, R. T. Rood, D. N. Schramm, J. W. Truran, and E. Vangioni-Flam, Preprint FERMILAB-Pub-338-A (September 1994), submitted to *Astrophys. J.*
- [22] A. Vidal-Majar, *Adv. Space Res.* **11**, 97 (1991); R. Ferlet, in *Astrochemistry of Cosmic Phenomena*, ed. P. D. Singh (Dordrecht: Kluwer), 85 (1992); J. L. Linsky *et al.*, *Astrophys. J.* **402**, 694 (1992).
- [23] T. M. Bania, R. T. Rood, and T. L. Wilson, in *Origin and Evolution of the Elements*, ed. N. Prantzos, E. Vangioni-Flam, and M. Casse (Cambridge: Cambridge Univ. Press), 107 (1993).
- [24] R. I. Epstein, J. M. Lattimer, and D. N. Schramm, *Nature* **263**, 198 (1976).
- [25] A. N. Gorbunov and V. M. Spiridonov, *Sov. Phys. JETP* **34(7)**, N.4, 600 (1958); Yu. M. Arkatov *et al.*, *Sov. J. Nucl. Phys.* **19**, 598 (1974).
- [26] Note that this is *not* evident from the figures and tables given in reference [11]. N. Yu. Gnedin, private communication.
- [27] for a review see R. Blandford and D. Eichler, *Phys. Rep.* **154**, 1 (1987).
- [28] T. K. Gaisser, *Cosmic Rays and Particle Physics* (Cambridge: Cambridge University Press, 1990).
- [29] A. M. Hillas, *Ann. Rev. Astron. Astrophys.* **22**, 425 (1984).

- [30] P. L. Biermann and P. A. Strittmatter, *Astrophys. J.* **322**, 643 (1987).
- [31] *Astrophysical Aspects of the Most Energetic Cosmic Rays*, eds.: M. Nagano and F. Takahara (World Scientific, Singapore, 1991).
- [32] C. T. Cesarsky, *Nucl. Phys. B* (Proc. Suppl.) **28B**, 51 (1992).
- [33] D. J. Bird *et al.*, *Phys. Rev. Lett.* **71**, 3401 (1993); *Astrophys. J.* **424**, 491 (1994).
- [34] D. J. Bird *et al.*, *Astrophys. J.* **441**, 144 (1995).
- [35] *Proc. of Tokyo Workshop on Techniques for the Study of Extremely High Energy Cosmic Rays*. (ICRR, Univ. of Tokyo, 1993).
- [36] S. Yoshida *et al.*, Institute of Cosmic Ray Research (Univ. of Tokyo) Report *ICRR-Report-325-94-20*, to appear in *Astropart. Phys.* (1995).
- [37] N. Hayashida *et al.*, *Phys. Rev. Lett.* **73**, 3491 (1994).
- [38] N. N. Efimov *et al.* in Ref. [31], p. 20.
- [39] T. A. Egorov in Ref. [35], p. 35.
- [40] P. Sommers in Ref. [35], p. 23.
- [41] G. Sigl, D. N. Schramm, and P. Bhattacharjee, *Astropart. Phys.* **2**, 401 (1994).
- [42] J. W. Elbert and P. Sommers, *Astrophys. J.* **441**, 151 (1995).
- [43] P. Bhattacharjee in Ref. [31].
- [44] P. Bhattacharjee, C. T. Hill, and D. N. Schramm, *Phys. Rev. Lett.* **69**, 567 (1992).
- [45] P. Bhattacharjee and N. C. Rana, *Phys. Lett.* **B246**, 365 (1990).
- [46] P. Bhattacharjee and G. Sigl, Preprint FERMILAB-Pub-94/387-A, to appear in *Phys. Rev. D*.
- [47] E. W. Kolb and M. S. Turner, *The Early Universe* (Addison-Wesley, New York, 1990).
- [48] B. Allen and E. P. S. Shellard, *Phys. Rev. Lett.* **64**, 119 (1990).
- [49] A. Albrecht and N. Turok, *Phys. Rev.* **D40**, 973 (1989).
- [50] D. P. Bennet and F. R. Bouchet, *Phys. Rev. Lett.* **60**, 257 (1988); **63**, 2776 (1989).
- [51] B. Allen and R. R. Caldwell, *Phys. Rev.* **D45**, 3447 (1992).
- [52] D. Austin, E. J. Copeland, and T. W. B. Kibble, *Phys. Rev.* **D48**, 5594 (1993).

- [53] V. S. Berezinsky and H. R. Rubinstein, *Nucl. Phys.* **B323**, 95 (1989).
- [54] R. Plaga, *Astrophys. J.* **424**, L9 (1994).
- [55] A. J. Gill and T. W. B. Kibble, *Phys. Rev.* **D50**, 3660 (1994).
- [56] X. A. Siemens and T. W. B. Kibble, Preprint *hep-ph/9412216*.
- [57] F. A. Aharonian, P. Bhattacharjee, and D. N. Schramm, *Phys. Rev.* **D46**, 4188 (1992).
- [58] W. Hu and J. S. Silk, *Phys. Rev.* **D48**, 485 (1993); see also references therein.
- [59] Ya. B. Zel'dovich and R. A. Sunyaev, *Astrophys. Space Sci.* **4**, 301 (1969).
- [60] R. Svensson and A. A. Zdziarski, *Astrophys. J.* **349**, 415 (1990).
- [61] C. E. Fichtel *et al.*, *Astrophys. J.* **217**, L9 (1977); **222**, 833 (1978); D. J. Thompson, and C. E. Fichtel, *Astron. Astrophys.* **109**, 352 (1982).
- [62] A. Vilenkin, *Phys. Rep.* **121**, 263 (1985).
- [63] V. S. Berezinsky, S. V. Bulanov, V. A. Dogiel, V. L. Ginzburg, and V. S. Ptuskin, *Astrophysics of Cosmic Rays* (North-Holland, 1990).
- [64] X. Chi, C. Dahanayake, J. Wdowczyk, and A. W. Wolfendale, *Astropart. Phys.* **1**, 239 (1993).

## Figure Captions

**Figure 1:** Maximal energy release in units of the CMBR energy density allowed by the constraints from the observed  $\gamma$ -ray background at 200MeV (dotted curve), CMBR distortions (dashed curve, from Ref. [13]), and  ${}^4\text{He}$ -photodisintegration as a function of redshift  $z$ . These bounds apply for instantaneous energy release at the specified redshift epoch. The logarithm is to the basis 10.

