

## Electroweak Jet Cascading in the Decay of Superheavy Particles

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We study decays of superheavy particles  $X$  into leptons. We show that they initiate cascades similar to QCD parton jets, if  $m_X \gtrsim 10^6$  GeV. Electroweak cascading is studied and the energy spectra of the produced leptons are calculated in the framework of a broken SU(2) model of weak interactions. As application, important for the Z-burst model for ultrahigh energy cosmic rays, we consider decays of superheavy particles coupled on tree level only to neutrinos and derive a stringent limit for these decays from the observed diffuse extragalactic  $\gamma$ -ray flux.

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A characteristic feature of high energy deep-inelastic scattering,  $e^+e^-$  annihilation, and decays of superheavy particles is the cascading of QCD partons. Despite the smallness of the QCD coupling  $\alpha_3$ , cascading occurs because the probability of parton splitting is enhanced by large logarithms for soft or collinear parton emission. Recently, similar logarithms were found to dominate also the EW radiative corrections at the TeV scale and above [1]. Evolution equations for these corrections, similar to the Dokshitzer-Gribov-Lipatov-Altarelli-Parisi (DGLAP) equations [2,3] in QCD but valid for a spontaneously broken theory, were derived in [4]. In this Letter, we want to draw attention to decays of particles much heavier than the EW scale: if these particles decay or annihilate, then the effects discussed above result in particle cascades which proceed through all  $U(1) \times SU(2) \times SU(3)$  interactions with couplings  $\alpha_1$ ,  $\alpha_2$ , and  $\alpha_3$  enhanced by logarithmic terms. While particular attention is given to EW lepton-boson (EWLB) cascades, our results could be generalized to an arbitrary spontaneously broken gauge group [5].

We shall illustrate the development of EW cascades by the example of a superheavy  $X$  particle with mass  $m_X \lesssim m_{\text{GUT}}$  decaying into leptons, which has an interesting physical application to the Z-burst model (see below). Let us discuss first the case when the  $X$  particle has on tree level only couplings to lepton pairs  $\bar{l}l$  and consider the decay mode  $X \rightarrow \bar{\nu}\nu$ . For  $m_X \gg m_Z$ , the mass of the Z boson is negligible compared to the available momentum flow  $Q^2 \leq m_X^2/4$ . Then, soft and collinear singularities generate large logarithms  $\ln^2(m_X^2/m_Z^2)$ , which can compensate the smallness of  $\alpha_2/(2\pi)$ .

To show the existence of the cascade, we consider the ratio  $R = \Gamma(X \rightarrow \bar{\nu}\nu Z)/\Gamma(X \rightarrow \bar{\nu}\nu)$ . Neglecting terms finite in the limit  $m_Z \rightarrow 0$ , it is given by

$$R = \frac{\alpha_2}{8\pi c_W^2} (\ln^2 \varepsilon + 3 \ln \varepsilon + \dots), \quad (1)$$

where  $\varepsilon = (m_Z/m_X)^2$  and  $c_W^2 = \cos^2 \vartheta_W$ . The  $\ln^2 \varepsilon$  term — which compensates the small coupling — is typical for soft and collinear singularities in the emission of vector particles [6]. For  $m_X \sim 10^6$  GeV, the decay probability into  $\bar{\nu}\nu$  and  $Z\bar{\nu}\nu$  becomes comparable,  $R \sim 0.5$ , signaling the breakdown of perturbation theory: the decay of a particle with mass  $m_X \gg 10^6$  GeV, even if coupled only to leptons, initiates a cascade, very similar to that known in QCD.

For the quantitative study of EWLB cascades, we consider a broken SU(2) gauge theory as a simplified model for the electroweak sector of the standard model (SM). Its particle content is a triplet of gauge bosons  $W_\mu^A$  with mass  $m_W$ , a physical Higgs  $h$  with mass  $m_h$ , and  $n_g = 3$  generations of leptons  $f = (l_L, \nu_{lL}), l_R$ , with  $l = e, \mu, \tau$ . The SU(2) doublet  $f$  is coupled to the gauge sector as  $g_2 t_{bc}^A \bar{f}_b \gamma^\mu \frac{1}{2}(1 - \gamma^5) f_c W_\mu^A$ .

Since the terms generating the logarithms in Eq. (1) are associated with collinear and/or soft emission of additional particles, we introduce splitting functions  $P_{ijk}(z)$ , viz.

$$\frac{d\Gamma_{n+1}}{d\Gamma_n} \approx d\omega_{i \rightarrow jk} = \frac{\alpha_2}{2\pi} \frac{dt}{t} dz P_{ijk}(z), \quad (2)$$

where  $d\Gamma_n$  is the decay width into  $n$  particles with virtuality  $t$ ,  $z$  is the energy fraction of the additional particle, and  $d\omega_{i \rightarrow jk}$  approximates  $d\Gamma_{n+1}/d\Gamma_n$  in the high-energy, small-angle scattering limit.

In the following, it is useful to distinguish between transverse and longitudinal modes of the gauge bosons,  $W_T$  and  $W_L$ . Going from QCD to our broken SU(2) model, there are only trivial changes for the transverse modes due to the different structure constants and the V-A coupling,

$$P_{W_T W_T W_T} = 2 \left[ \frac{1-z}{z} + \frac{z}{1-z} + z(1-z) \right], \quad (3)$$

$$P_{W_T ff} = \frac{n_g}{4} [z^2 + (1-z)^2], \quad (4)$$

$$P_{ffW_T} = \frac{3}{2} P_{ff'W_T^\pm} = 3P_{ffW_T^0} = \frac{3}{4} \frac{1+z^2}{1-z}. \quad (5)$$

External longitudinal modes  $W_L$  can be replaced by the corresponding Goldstone bosons  $\chi_W$  in  $S$ -matrix elements for  $s \gg m_W^2$  [7]. Since  $m_W$  acts as an infrared cutoff and is assumed to be well above the experimental resolution, virtual corrections and real bremsstrahlung do not mix. Replacing  $W_L$  by the scalars  $\chi_W$  makes clear that the splitting functions involving  $W_L$  do not contain terms singular in  $z$ . The longitudinal modes  $W_L$  are therefore subdominant compared to  $W_T$  and we will neglect them.

Finally, we discuss whether the Higgs participates actively in the cascade evolution. The Higgs can split into pairs of fermions or gauge bosons. In the first case, all Yukawa couplings except the one to the top quark  $t$  can be neglected; moreover, collinear singularities are absent. Therefore, only the processes  $t \rightarrow th$  and  $h \rightarrow tt$  have a minor importance at  $Q^2$  close to the GUT scale. In the second case, i.e., when splittings involve Higgses and gauge bosons, we cannot define splitting functions in the conventional sense. Instead,  $d\omega_{h \rightarrow W_L W_L}$  and  $d\omega_{h \rightarrow W_T W_T}$  are proportional to  $g^2/t^2$ . The resulting branching probability is negligible.

Hence, the EW cascade is supported essentially by the same splitting modes as in the unbroken theory. The proof of the factorization of both collinear and soft singularities in the EW theory follows the same logic as in QCD [2,3] and the effects of the spontaneous symmetry breaking result at  $s \gg m_W^2$  only in subleading corrections [8]. Thus one obtains the EW evolution equations in the Altarelli-Parisi form [4] and is able to establish the SU(2) charge coherence picture in the usual way [9]. The evolution equations in integral form including coherence effects are

$$\begin{aligned} D_{i \rightarrow j}(Q^2, x, t) &= \Delta_s^i(Q^2, t) \delta_i^j \delta(1-x) \\ &+ \sum_{kl} \int_t^{Q^2} \frac{dt'}{t'} \int_x^{1-\varepsilon_l(t')} \frac{dz}{2\pi z} \Delta_s^i(Q^2, t') \\ &\times \alpha_2 [z^2(1-z)^2 t'] P_{ikl}(z) D_{k \rightarrow j}(x/z, z^2 t'), \end{aligned} \quad (6)$$

where  $D_{i \rightarrow j}(Q^2, x, t)$  defines the energy distribution of particles  $j$  with energy fraction  $x$  at scale  $t$  produced by the parent particle  $i$  at scale  $Q^2$ . Here,  $\varepsilon_l(t') \approx m_l/\sqrt{t'}$ , where  $m_l$  is the mass of the particle  $l$ ,  $t = q^2/[z(1-z)]$ ,  $q^2$  the virtuality of the parent particle, and  $z$  the energy fraction of the produced one. We use  $\alpha_2(t) = \alpha_2(m_Z^2)/[1 + (b/4\pi) \ln(t/m_Z^2) \alpha_2(m_Z^2)]$  with  $b = 19/6$  for one Higgs doublet and  $\alpha_2(m_Z^2) \sim 1/29.6$ .

In Eq. (6), we have introduced the Sudakov form factors  $\Delta_s^i(Q^2, t)$  giving the probability of no branching for

the parton of type  $i$  in the scale range between  $Q^2$  and  $t$ ,

$$\begin{aligned} \ln \Delta_s^i(Q^2, t) &= - \sum_{kl} \int_t^{Q^2} \frac{dt'}{t'} \int_{\varepsilon_k(t')}^{1-\varepsilon_l(t')} \\ &\times \frac{dz}{2\pi} \alpha_2 [z^2(1-z)^2 t'] P_{ikl}(z). \end{aligned} \quad (7)$$

Equations (6) and (7) are very similar to the analogous ones in QCD and allow use of the probabilistic scheme of Ref. [10]. There is, however, an important difference between the EWL and QCD cascades in the physical meaning of the cutoff value  $\varepsilon_k(t')$ . In QCD, it is defined by the scale  $q_{\min}^2$  at which the perturbative evolution of the cascade is terminated,  $\varepsilon_{\text{QCD}}(t') = \sqrt{q_{\min}^2/t'}$ . By contrast,  $q_{\min}^2$  of the EWL cascade is given by the physical masses of the particles produced in the splitting  $i \rightarrow kl$ . For  $W_T \rightarrow ff$ ,  $\varepsilon_k(t') \approx 0$ , i.e., it is much smaller than the one in  $W_T \rightarrow W_T W_T$ , where  $\varepsilon_k(t') \approx m_W/\sqrt{t'}$ . As a consequence, the  $W_T$  does not dominate the cascade evolution as strongly as gluons in QCD and the number of  $W_T$  leaving the cascade is smaller than the number of leptons.

The EWL cascade does not terminate abruptly at a certain  $Q^2$ . Instead, the probability for no further branching increases smoothly from  $\Delta_s^i(Q^2, 0) = 0$  for  $Q^2 \rightarrow \infty$  to  $\Delta_s^i(0, 0) = 1$ . Nevertheless, one can define  $Q_{\text{cr}}^2$  by  $\Delta_s^i(Q_{\text{cr}}^2, 0) = 0.5$ , i.e., by the virtuality  $Q_{\text{cr}}^2$  at which the probabilities for further branching and for nonbranching become equal. Then  $Q_{\text{cr}}$  is given by  $Q_{\text{cr}} \sim 10^6$  GeV and the probability for further splittings decreases indeed drastically for lower  $Q$ .

We choose to solve the evolution equation by means of a Monte Carlo simulation. This method gives us the advantage of including non-Abelian charge coherence effects through angular ordering. In Fig. 1, we show the spectra  $dN/dl$ , where  $l = \ln(1/x)$  and  $x = 2E/m_X$ , of leptons and  $W$ 's for  $m_X = 10^{10}$ ,  $10^{13}$ , and  $10^{16}$  GeV and

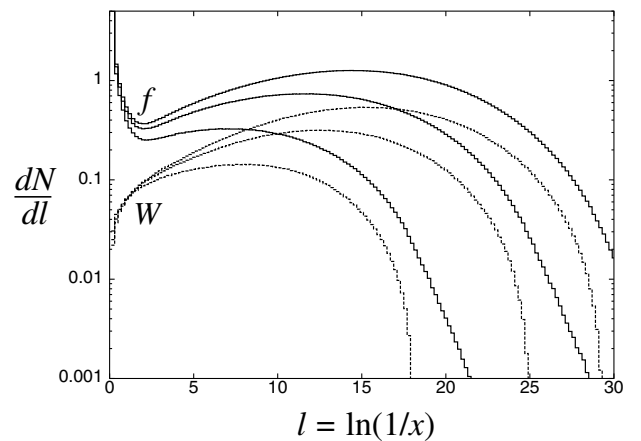


FIG. 1. Energy spectra  $dN/dl$  of leptons  $f$  (solid lines) and  $W$ 's (broken lines) produced in  $X \rightarrow \bar{\nu}\nu \rightarrow \text{all}$ ; for  $m_X = 10^{16}$ ,  $10^{13}$ , and  $10^{10}$  GeV from top to bottom.

for the decay mode  $X \rightarrow \bar{\nu}\nu \rightarrow$  all. The average number of produced leptons and  $W$ 's increases from 9 for  $m_X = 10^{10}$  GeV to 47 for  $m_X = 10^{16}$  GeV, where we consider the  $\mu^\pm$ ,  $\tau^\pm$ , and the  $W^a$  as stable particles. The first bin contains between 7% ( $m_X = 10^{10}$  GeV) and 0.2% ( $m_X = 10^{16}$  GeV) of prompt neutrinos with  $E = m_X/2$  from noncascading decays. Leptons which stop branching after the first splitting produce a tail superimposed to the usual Gaussian, clearly visible for  $l \lesssim 3$ .

Let us discuss now the inclusion of QCD partons in our model. Since the EW gauge bosons split also into quarks  $q$ , there will be a mutual transmutation of "leptonic" and "QCD" cascades. The shape of the hadron energy spectra is, however, only marginally influenced by the leptons, first because the QCD cascade is determined mainly by gluons  $g$  and second because the probability of  $q \rightarrow q + g$  is much larger than of  $q \rightarrow q + W$ . On the other hand, the splittings  $W_T \rightarrow qq$  act as a sink for the particles and energy of the EWL cascade. These splittings are taken into account in the MC simulation and the spectra shown in Fig. 1.

We have calculated also (see Table I) a quantity interesting for decays of  $X$  particles in the Universe: the fraction of energy  $f_{\text{em}}$ , transferred to electrons and photons in the decay mode  $X \rightarrow \bar{\nu}\nu$  as a function of  $m_X$ . These particles initiate electromagnetic (e.m.) cascades in the intergalactic space, through interactions with microwave (CMB) photons. In the calculation of  $f_{\text{em}}$ , we assumed that all hadrons produced in  $W_T \rightarrow q\bar{q} \rightarrow$  hadrons are pions, while we used the SM branching ratios for  $\mu$  and  $\tau$  decays.

The above calculations have interesting consequences for ultrahigh energy cosmic rays (UHECR) produced by superheavy dark matter (DM) [11] and, in particular, by  $Z$  bursts [12]. In the latter model, UHECR are produced through the resonant production of  $Z$  bosons in the collisions of UHE neutrinos with DM neutrinos,  $\nu + \nu \rightarrow Z^0$ . The resonant energy of UHE neutrino is  $E_0 = m_Z^2/2m_\nu = 4.2 \times 10^{21} m_{\text{eV}}^{-1}$  eV, where  $m_{\text{eV}}$  is the mass of the DM neutrino in eV.

Decays of these  $Z$  bosons produce UHECR which energy spectrum, according to recent calculations [13,14], has a weak Greisen-Zatsepin-Kuzmin cutoff and explains well the observations. A remarkable feature of this model is that the production of UHE particles with energies higher than  $10^{20}$  eV does not involve exotic elementary particle physics, while its drawback consists in the necessity of an enormous flux of resonant neutrinos.

TABLE I. Energy fraction  $f_{\text{em}}$  transferred to electrons and photons in the decay  $X \rightarrow \bar{\nu}\nu$  of a particle with mass  $m_X$ .

| $m_X/\text{GeV}$   | $10^6$ | $10^8$ | $10^{10}$ | $10^{12}$ | $10^{14}$ | $10^{16}$ |
|--------------------|--------|--------|-----------|-----------|-----------|-----------|
| $f_{\text{em}}/\%$ | 11     | 15     | 17        | 19        | 20        | 20        |

A widely discussed possibility is the production of UHE neutrinos in astrophysical sources. This is an unrealistic option because it requires too high luminosities of the astrophysical sources: we can estimate the required flux of resonant neutrinos  $I_\nu(E_0)$  and therefrom the neutrino energy density  $\omega_\nu$  as  $\omega_\nu \approx (2.4 - 3.6) \times 10^{-13} m_{\text{eV}}^{-0.5} \text{erg/cm}^3$ . In these calculations, we have used that for the flat spectra generated by heavy particle decays, such as  $Z$  bosons and  $X$  particles, UHE photons dominate at the highest energies  $E \geq 1 \times 10^{20}$  eV [15]. The resulting neutrino luminosity of a source, estimated as  $L_\nu \sim \omega_\nu/(n_s t_0)$ , where  $n_s$  is the source density and  $t_0$  the age of the Universe, is too high:  $(8 - 12) \times 10^{44}$  erg/s, if the sources are normal galaxies, and  $(8 - 12) \times 10^{46}$  erg/s in the case of Seyfert galaxies. Our result confirms similar conclusions of Ref. [14], reached on the basis of a limit due to the diffuse  $\gamma$ -ray background.

Alternatively,  $Z$ -burst neutrinos can be generated by decays of superheavy particles, either existing as DM particles or generated by topological defects. Usually, this possibility is ruled out by the diffuse  $\gamma$ -ray background [16]. However, Ref. [17] suggested that this limit can be evaded, if  $X$  particles decay exclusively to neutrinos.

We shall demonstrate now that electroweak cascading rules out this last hope, too.

EW cascading modifies the e.m. cascading upper limit for UHE neutrinos [16,18] by the factor  $f_\nu/f_{\text{em}}$ , where  $f_\nu$  is the fraction of energy  $m_X$  transferred to neutrinos. Then the upper limit on the UHE neutrino flux is given as

$$I_\nu(E) < \frac{c}{4\pi} \frac{f_\nu}{f_{\text{em}}} \frac{\omega_{\text{cas}}}{E^2}, \quad (8)$$

with  $\omega_{\text{cas}} \approx 2 \times 10^{-6}$  eV/cm<sup>3</sup> according to Energetic Gamma Ray Experiment Telescope observations. Taking  $I_\nu(E_0)$  from Eq. (8), one can calculate the flux of UHE photons produced in  $Z$  bursts. At  $E \geq 1 \times 10^{20}$  eV, this flux should be of the order of the observed UHECR flux. With  $f_{\text{em}} = 0.2$  from Table I, we obtain  $E^3 I_\gamma(E) \approx (5 - 8) \times 10^{21} m_{\text{eV}} \text{eV}^2/\text{m}^2\text{s sr}$ , which is almost 3 orders of magnitude less than observed.

In conclusion, EW cascading in decays of superheavy particles results in EWL cascades similar to the cascade of QCD partons, if  $m_X > 10^6$  GeV. Thereby, the flux of prompt neutrinos, e.g., in annihilations of DM particles is reduced. EW cascades allow a probabilistic interpretation and exhibit destructive coherence at small  $x$ . The generation of electrons and photons, which are able to start e.m. cascades on the CMB, in the EW cascade allows one to exclude those  $Z$ -burst models, in which the resonant neutrinos are produced in  $X \rightarrow \bar{\nu}\nu$  decays. Thus, the production of resonant neutrinos in astrophysical sources and by decays of superheavy particles (both as DM particles or produced by topological defects) result in neutrino fluxes too low for the  $Z$ -burst model. The only possibility left is the oscillation of sterile neutrinos, e.g., in hidden-sector/mirror models [16], into ordinary ones.

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