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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 (EW) 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 stringent limit for these decays from the observed diffuse extragalactic γ-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 α_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 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)×SU(2)×SU(3) interactions with couplings α_1, α_2 and α_3 enhanced by logarithmic terms. Particular attention is given to EW Lepton-Boson (EWLB) cascades developing without QCD partons.

We shall illustrate the development of EW cascades by the example of a superheavy X-particle with mass $m_X \lesssim m_{\rm GUT}$ 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 $X \to \bar{\nu}\nu$. 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 $R = \Gamma(X \to \bar{\nu}\nu Z)/\Gamma(X \to \bar{\nu}\nu)$. Neglecting terms finite in the limit $m_Z \to 0$, it is given by

$$R = \frac{\alpha_2}{8\pi c_W^2}\left(\ln^2\varepsilon + 3\ln\varepsilon + \ldots\right), \qquad (1)$$

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 [5]. 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 break-down 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_μ^A with mass m_W, a physical Higgs h with mass m_h and $n_g = 3$ generations of leptons $f = (l_L, \nu_{l,L})$, l_R, with $l = e$, μ, τ. The SU(2) doublet f is coupled to the gauge sector as $g_2 t^A_{bc}\bar{f}_b\gamma^\mu\frac{1}{2}(1-\gamma^5)f_c W_\mu^A$.

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.

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

where $d\Gamma_n$ is the decay width into n particles with virtuality t, z is the energy fraction of

the additional particle, and $d\omega_{i\to jk}$ 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,

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

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

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

External longitudinal modes W_L can be replaced by the corresponding Goldstone bosons χ_W in S-matrix elements for $s \gg m_W^2$ [6]. Since m_W acts as an infrared cut-off and is assumed to be well above the experimental resolution, virtual corrections and real bremsstrahlung do not mix. Replacing W_L by the scalars χ_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 $t \to th$ and $h \to tt$ 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, $d\omega_{h\to W_LW_L}$ and $d\omega_{h\to W_TW_T}$ 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 [7]. 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 [8]. The evolution equations in integral form including coherence effects are

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

where $D_{i\to j}(Q^2,x,t)$ 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,

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

Equations (6-7) are very similar to the analogous ones in QCD and allow to use the probabilistic scheme of Ref. [9]. There is however an important difference between the EWLB and QCD cascades in the physical meaning of the cut-off value $\varepsilon_k(t')$. In QCD, it is defined by the scale $q^2_{\rm min}$ at which the perturbative evolution of the cascade is terminated, $\varepsilon_{\rm QCD}(t') = \sqrt{q^2_{\rm min}/t'}$. By contrast, $q^2_{\rm min}$ of the EWLB cascade is given by the physical masses of the particles produced in the splitting $i \to kl$. For $W_T \to ff$, $\varepsilon_k(t') \approx 0$, i.e., it is much smaller than the one in $W_T \to W_TW_T$, 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 EWLB cascade does not terminate abruptly at a certain Q^2. Instead, the probability for no further branching increases smoothly from $\Delta_s^l(Q^2,0) = 0$ for $Q^2 \to \infty$ to $\Delta_s^l(0,0) = 1$. Nevertheless, one can define $Q^2_{\rm cr}$ by $\Delta_s^l(Q_c^2,0) = 0.5$, i.e. by that virtuality $Q^2_{\rm cr}$ at which the probabilities for further branching and for non-branching become equal. Then $Q_{\rm cr}$ is given by $Q_{\rm cr} \sim 10^6$ GeV and the probability for further splittings decreases indeed drastically for lower Q.

We choose to solve the evolution equation with 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 for the decay mode $X \to \bar{\nu}\nu \to$ 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 2‰ ($m_X = 10^{16}$ GeV) of prompt neutrinos with $E = m_X/2$ from no-cascading decays. Leptons which stop branching after the first splitting produce a tail superimposed to the usual Gaussian, clearly visible for $l \lesssim 3$.

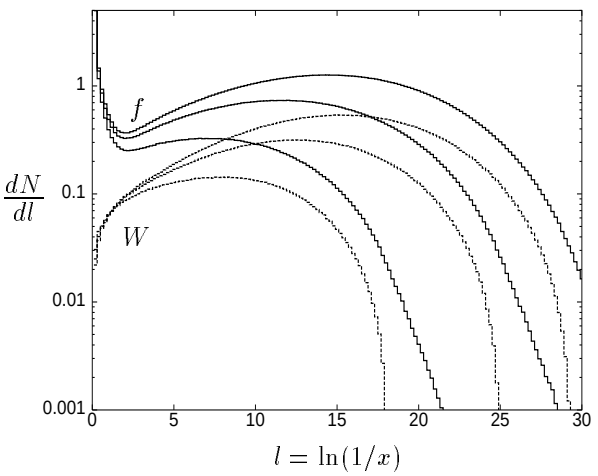

FIG. 1. Energy spectra dN/dl of leptons f (solid lines) and W's (broken lines) produced in $X \to \bar{\nu}\nu \to$ all; for $m_X = 10^{16}$, 10^{13} and 10^{10} GeV from top to down.

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 $q \to q + g$ is much larger than of $q \to q + W$. On the other hand, the splittings $W_T \to qq$ act continuously as a sink for the particles and energy of the EWLB cascade. These splittings are taken into account in the MC simulation and the spectra shown in Fig. 1.

We have calculated also (see Table 1) a quantity interesting for decays of X-particles in the Universe: the fraction of energy $f_{\rm em}$, transferred to electrons and photons in the decay mode $X \to \bar{\nu}\nu$ as 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 $f_{\rm em}$, we assumed that all hadrons produced in $W_T \to q\bar{q} \to$ hadrons are pions, while we used the SM branching ratios for μ and τ decays.

The above calculations have interesting consequences for Ultra High Energy Cosmic Rays (UHECR) produced by superheavy Dark Matter (DM) [10] and, in particular, by Z-bursts [11]. In the latter model, UHECR are produced through the resonant production of Z-bosons in the collisions of UHE neutrinos with DM neutrinos, $\nu + \bar{\nu} \to Z^0$. The resonant energy of UHE neutrino is $E_0 = m_Z^2/2m_\nu = 4.2 \times 10^{21} m_{\rm eV}^{-1}$ eV, where $m_{\rm eV}$ is the mass of the DM neutrino in eV.

Decays of these Z-bosons produce UHECR which energy spectrum, according to recent calculations [12,13], 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.

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 ω_ν as $\omega_\nu \approx (2.4 - 3.6) \times 10^{-13} m_{\rm eV}^{-0.5}$ 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. [13], reached on the basis of a limit due to the diffuse γ-ray background.

Alternatively, Z-burst neutrinos can be generated by decays of superheavy particles, either existing as DM particles or generated by topological defects. In the usual models this possibility is ruled out by the diffuse γ-ray background [16]. However, in Ref. [14] it is suggested that this limit can be evaded, if X-particles decay exclusively to neutrinos, i.e. $X \to \bar{\nu}\nu$.

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 [16,17] by the factor $f_\nu/f_{\rm em}$, where f_ν is the fraction of energy m_X transferred to neutrinos. Then the upper limit on the UHE neutrino flux is given as

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

with $\omega_{\rm cas} \approx 2\times 10^{-6}$ eV/cm^3 according to EGRET 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 $f_{\rm em} = 0.2$ from Table 1, we obtain $E^3 I_\gamma(E) \approx (5-8)\times 10^{21} m_{\rm eV}$ eV^2/m^{2}s sr, which is almost three orders of magnitude less than observed.

Summary—EW cascading in decays of superheavy particles results in EWLB 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 to exclude those Z-burst models, in which the resonant neutrinos are produced in $X \to \bar\nu\nu$ 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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[1] P. Ciafaloni and D. Comelli, Phys. Lett. B **446**, 278 (1999); M. Beccaria *et al.*, Phys. Rev. D **61**, 073005 (2000); V. S. Fadin *et al.*, Phys. Rev. D **61**, 094002 (2000); W. Beenakker and A. Werthenbach, Phys. Lett. B **489**, 148 (2000); M. Hori, H. Kawamura and J. Kodaira, Phys. Lett. B **491**, 275 (2000); A. Denner and S. Pozzorini, Eur. Phys. J. C **18**, 461 (2001); J. H. Kühn *et al.*, Nucl. Phys. B **616**, 286 (2001). For early works see W. Beenakker *et al.*, Nucl. Phys. **B410**, 245 (1993), Phys. Lett. **B317**, 622 (1993). For a recent review see M. Melles, hep-ph/0104232.

[2] V. N. Gribov and L. N. Lipatov, Sov. J. Nucl. Phys. **15**, 438 and 675 (1972).

[3] G. Altarelli and G. Parisi, Nucl. Phys. **B126**, 298 (1977); Yu. L. Dokshitzer, Sov. Phys. JETP **46**, 641 (1977).

[4] M. Ciafaloni, P. Ciafaloni, D. Comelli, Phys. Rev. Lett. **88**, 102001 (2002).

[5] S. Weinberg, Phys. Rev. **140**, B516 (1965); J. R. Ellis, M. K. Gaillard and G. G. Ross, Nucl. Phys. B **111**, 253 (1976).

[6] J. M. Cornwall, D. N. Levine and G. Tiktopoulos, Phys. Rev. **D10**, 1145 (1974); B. W. Lee, C. Quigg and H. B. Thacker, Phys. Rev. D **16**, 1519 (1977); M. S. Chanowitz and M. K. Gaillard, Nucl. Phys. B **261**, 379 (1985).

[7] S. Pozzorini, hep-ph/0201077 and references therein.

[8] A. Bassetto, M. Ciafaloni and G. Marchesini, Phys. Rept. **100**, 201 (1983); Yu. L. Dokshitzer *et al.*, *Basics of Perturbative QCD*, Editions Frontières 1991.

[9] G. Marchesini and B. R. Webber, Nucl. Phys. **B238**, 1 (1984).

[10] V. Berezinsky, M. Kachelrieß, and A. Vilenkin, Phys. Rev. Lett. **79**, 4302 (1997); V. A. Kuzmin and V. A. Rubakov, Phys. Atom. Nucl. **61**, 1028 (1998).

[11] D. Fargion, B. Mele and A. Salis, Astrophys. J. **517**, 725 (1999); T.J. Weiler, Astropart. Phys. **11**, 303 (1999).

[12] Z. Fodor, S. D. Katz and A. Ringwald, hep-ph/0105064 and hep-ph/0203198. See also S. Yoshida, G. Sigl and S. J. Lee, Phys. Rev. Lett. **81**, 5505 (1998).

[13] O. E. Kalashev, V. A. Kuzmin, D. V. Semikoz and G. Sigl, hep-ph/0112351.

[14] G. Gelmini and A. Kusenko, Phys. Rev. Lett. **84**, 1378 (2000).

[15] V. Berezinsky, P. Blasi and A. Vilenkin, Phys. Rev. D **58**, 103515 (1998).

[16] V. S. Berezinsky and A. Vilenkin, Phys. Rev. D **62**, 083512 (2000).

[17] V. S. Berezinsky, in Proc. of "Neutrino-77", Baksan , USSR, ed. M.A.Markov **1**, 177 (1977).

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

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