Kaluza-Klein Contamination in Fermi Accelerated Environments

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Abstract

Astrophysical constraints of new physics are often limited to weakly interacting light particles, such as axions, the Kaluza-Klein (KK) gravitons from the ADD model, sterile neutrinos and unparticles. We discuss the possibility for an astrophysical scenario to (dis)confirm new physics for heavy particles beyond TeV energy scale. In our scenario, the KK protons (the KK excited quarks/gluons within protons) within the framework of universal extra dimensions (UEDs), are produced by high energy $p + p$ collisions in Fermi accelerated environments, with protonic isotropic spectrum $dN/dE \propto E^{-2}$ up to at least $10^{18}$ eV. Thus, because they are also electrically charged, they should be re-accelerated by mechanism similar to normal protons. The KK states (no matter whether they have already decayed to the lightest KK particle or not) should contaminate $10^{-5}$ to $10^{-2}$ of cosmic-ray events for some fixed energy $E$ (within some suitable assumptions). Hence, if we have techniques to identify them from air shower data, we can constrain UEDs scenario. Our method is an “existence proof” that we can constrain new physics beyond TeV scale or much higher by classical astrophysical scenarios, which can also be generalized to supersymmetric models, the bulk Standard Model fields within the RS model, and the endlessly emerging new models. Moreover, it can exploit domains which have no possibility to be studied in terrestrial experiments.

Key words: cosmic-rays; Fermi mechanism; Kaluza-Klein states; models beyond the Standard Model; universal extra dimensions
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1 Introduction

Brane-world scenarios, such as the Arkani-Hamed-Dimopoulos-Dvali (ADD) model [1, 2, 3] and Randall-Sundrum (RS) model [4, 5], give an alternative framework for solving the hierarchy problem. Within the ADD model, the Standard Model (SM) fields are confined to a 3-brane ((3 + 1)-dimensional spacetime) while gravitons propagate freely in a torus compactified bulk space (large extra dimensions), and the gravitational coupling constant $G = 1/M_{\text{pl}}^2$ observed in our (3 + 1)-dimensional world is an effective one. So it is natural to understand why $G$ is so small. The RS model solves the same problem by a slice of $AdS_5$ spacetime. Some superstring-inspired descriptions make these models more attractive.

Some lineage scenarios of the ADD model also allow the SM particles propagating, to some extent, in the bulk spaces. Beside gravitons, such kind of SM particles can also be Kaluza-Klein (KK) excited in the extra dimensions, thus give abundant physical phenomena. A natural extension of the ADD model is let the brane with a finite thickness and complex substructures [6]. An example is given in Ref. [7]. In this scenario, quarks and leptons are confined to different branes, while the Higgs and SM gauge fields are sandwiched in, hence the possibility for proton decay can be exponentially suppressed. The KK excited gauge particles (which may be called branons) can be either baryophobic or lepto-phobic, because they feel nontrivial on the brane substructures. Another natural extension is asymmetrical compactification, which has two (as minimum, maybe more) separate compactification scales. The “very large” extra dimensions let only graviton propagate, just as what the ADD model says, but the $\text{TeV}^{-1}$ scale “large” extra dimensions, may have the SM fields extending [3]. Of course we can on the other hand keep all compact dimensions at $\text{TeV}^{-1}$ scale rather than make the compactification asymmetric. However, the advantage of solving the hierarchy problem in the ADD model is bereft.

In the case of that the SM fields also extend to some extra dimensions, to obtain chiral fermions in the 4-dimensional effective theory, we have only two ways to go: (i) to confine fermions to branes only [9], or (ii) to impose bulk fermions orbifold boundary conditions [10]. Universal extra dimensions (UEDs) [11] scenario is an example of the second approach. In UEDs scenario, all SM fields can propagate in these “universal” extra dimensions, and conservation of momentum in the universal dimensions turns to conservation of the KK number in our (3 + 1)-dimensional world. For two or more universal extra dimensions, the naïve KK mode sums diverge when the KK tower number $N_{KK} \to \infty$. So, let us consider only one universal extra dimension as minimal universal extra dimensions (MUEDs) in this context. In this case, $S^1/\mathbb{Z}_2$ orbifold compactification is assumed, and the KK mass eigenvalues have a simple form $M_{n}^{KK} = n/R$, where $R$ is the compactification scale. For the reason that only loop diagrams can contribute electroweak observables by the KK number conservation, the experimental bound for UEDs is only $M_{1}^{KK} = 1/R \geq 300 \text{GeV}$. In the tree level, the mass spectrum of the KK excited SM particles has the form $M_{SM,n} = \sqrt{(M_{n}^{KK})^2 + M_{SM}^2}$, where $M_{SM}$ is the zero-mode on-shell mass of the corresponding SM particles. Hence, their masses are level-by-level highly degenerated when $M_{SM} \ll M_{n}^{KK}$. However, when radiative corrections are concerned [12, 13], the mass degeneration is broken, to some extend, as
$M_{g,n} > M_{Q,n} > M_{q,n} > M_{W^\pm,n} \sim M_{Z,n} > M_{L,n} > M_{l,n} > M_{\gamma,n} \sim M_n$, where $g$ denotes gluon, $Q (L)$ denotes weak-doublet quark (lepton), $q (l)$ denotes weak-singlet quark (lepton), and $\gamma$ denotes photon. The KK number conservation breaks down to a KK parity that the even and odd KK numbers cannot transform to each other. The correction scale depends on some unknown parameters; however, some reasonable choice of parameters shows that the largest correction $\Delta M_{g,n}$ may be as large as 10% [13]. Hence, heavier KK excited states should cascade decay (by the KK conserving or even violating interactions) to the lightest KK particle (LKP) $\gamma_1$ which is stable [14], by emitting soft SM particles. When considering the possibilities of experimental discovery, one always assumes that the lifetimes of heavier KK excited states are sufficiently short, thus the states can decay within the collider; however, it is not supposed to do so. The total width cannot in fact be calculated by the failure of reconstructing the Breit-Wigner resonance [15]. Notice that the decay rate of an unstable particle $d\Gamma \propto 1/m_A$ in the phase space formula, the lifetime $\tau \propto m_A$ where $m_A$ is the mass of the decay particle. If the $\Delta M$s are smaller for a set of parameters different from in [13], or soft SM cascade processes are suppressed by other reasons, the lifetimes of heavier KK excited states should be even longer. Specific calculations for whether the not-very-short-lived KK excited states can affect other more mature scientific scenarios, such as disturb predictions of Big-bang nucleosynthesis, or distort the cosmic microwave background, are need; however, they are beyond our scope of our paper. Some naïve considerations show that all of them are not very crucial, because we do not really need longer-lived KK excited states (although the long-lived ones are also possibilities we shall consider in the identification section in §4), but some not-very-short-lived KK excited states to suffer the time scale of Fermi acceleration (which is maybe $\sim s$ or much shorter), which is much shorter than the time scale of the scenarios we mentioned above. So we assume that the lifetimes of heavier KK excited states are long enough to suffer the astrophysical scenario we draw in this paper.

Astrophysical constraints of new physics are often limited to weakly interacting light particles, such as axions [16, 17, 18, 19, 20, 21, 22, 23, 24, 25, 26, 27], the KK gravitons from the ADD model [3, 28, 29, 30, 31, 32, 33, 34], sterile neutrinos [35, 36, 37, 38] and unparticles [39]. We want to construct an astrophysical scenario to (dis)confirm new physics for heavy particles beyond TeV energy scale. Notice that in Fermi accelerated environments, protons in a power law spectrum up to at least $10^{18}$ eV should be produced (even if we have already derived an overall Lorentz factor $\Gamma \sim 300$), thus $p + p$ collisions up to a tremendously large energy should happen there, which we cannot even imagine in terrestrial experiments. However, we have to brain storm to know their happenings. In this paper, we construct a scenario which may (dis)confirm UEDs by high energy observation of cosmic-rays. This scenario may or may not have opportunities to give stronger bounds than colliders, because of the large uncertainties in our estimations, and the technical details of lots of synergic scientific domains (which we cannot discuss at length in this paper). However, it is at least an “existence proof” for this kind of methodology. It can also explore domains which have no possibility to be studied in terrestrial experiments. Some similar scenario in Ref. [40], also suggested the production of some kind of strongly interacting massive particles in $p + p$ collisions in astrophysical environments; however, our
scenario have a lot of advantages than theirs. The advantages rise mainly because (i) one of the protons in their scenario stays at rest, but both of the protons in our scenario are Fermi accelerated, and (ii) our KK excited states suffer an additional accelerated process. Detail comparisons are given in §6.

In our scenario, the KK protons (with either KK excited quarks or gluons in it) are produced by $p + p$ collisions in Fermi accelerated environments. Both the original Fermi mechanism or diffusive shock accelerating model have an isotropic spectrum $dN/dE \propto E^{-2}$ up to at least $10^{18}$ eV, hence they are okay for our purpose. The KK protons should also be accelerated just as normal protons by the same mechanism; however, they should have different properties than normal ones. Beside being discovered one by one from air shower data directly, they may be accelerated to energies normal protons cannot be accelerated to, or they (or their decayed final state) may contaminate significant amount of ultra-high-energy cosmic-ray (UHECR) events because of an overall energy shift, both of which may make them a discovery. In §2 we calculate the cross section and production rate of the KK protons. We show that the production rate may be large enough to make meaningful scientific constraints. In §3 we discuss the accelerated property of them in Fermi accelerated environments, making a comparison with normal protons. We notice that the KK states should contaminate $10^{-5}$ (for special sources) to $10^{-2}$ (for diffuse flux) of cosmic-ray events for some fixed energy $E$ (if assuming the optical depth $\tau_{pp} = 1$), which are not too small a sample to be discovered by air shower detection. We also discuss the propagating properties of them related to soft photon interactions. In §4 we consider the probabilities to identify them (or their decayed final state) from other cosmic-ray particles from air shower data. If it can be done so, our method can have larger possibilities to give smaller parameter space for UEDs than other methods. In §5 we calculate the possible constraints of the KK cosmic-ray flux from neutrino detectors; however, the constraints are very loose for current scientific equipments to affect our former estimations. We discuss our results and draw the possibilities to generalize our method to other new physics models in §6.

2 Producing of the KK Protons

The KK particles can be produced by high energy collisions, both in terrestrial experiments and astrophysical environments. We may want to consider a list of some not-very-exotic processes, such as $\gamma + \gamma, \ e + e, \ p + p, \ p + n, \ e + \gamma$ and $p + \gamma$. Collisions including electrons are very tempting; however, electrons always have limited energy because of synchrotron radiation. Ultra-high-energy (UHE) photons with energy $10^{12-14}$ eV or larger may be produced in astrophysical scenarios [11]; however, we leave behind this possibility elsewhere because of page limitations. Collisions including neutrons are very interesting; however, uncharged neutrons are always less energetic than protons, thus they make a less center-of-mass energy $\sqrt{s}$ for $p + n$ interactions than $p + p$ ones. In addition, we lack trustworthy parton distribution functions (PDFs) $f(x, Q)$ for high $Q$ neutrons. Hence, we will only discuss $p + p$ collisions in this paper.
2.1 \( p + p \rightarrow (\text{the KK states}) \) Cross Sections

The \( p + p \) cross sections to produce the KK bosons and fermions in the framework of UEDs have already been calculated in [42, 43]. However, their motivation is mainly on whether we have possibilities to confirm them on synchrotrons (especially Tevatron I, II and the Large Hadron Collider (LHC)), thus they focused their calculations on a center-of-mass energy \( \sqrt{s} \) of 1 TeV, 2 TeV and 14 TeV respectively and let \( M_{1}^{KK} \) as a free parameter. As a result, their calculations are not suitable for our purpose. So we redo the calculations for a set of parameters \( \sqrt{s} \). In this paper, we always fix \( M_{1}^{KK} = 1/R = 350 \text{ GeV} \) for simplicity.

In our calculation, we assume the KK states are sufficiently stable (see the discussions in §1). We use the amplitude-squared \( \sum|\mathcal{M}|^2 \) calculated by Macesanu et al. [43]. Their expression for total cross section is

\[
\sigma_{\text{tot}}^{KK} = \frac{1}{4\pi} \sum_{j} \sum_{n} \int_{\rho_{n}}^{1} dx_A \int_{\rho_{n}/x_A}^{1} dx_B \times f_1(x_A, Q) f_2(x_B, Q) \times \int_{0}^{\pi} \sin \theta d\theta \sum_{j} |\mathcal{M}_j|^2 \frac{1}{S!} \sqrt{\frac{1 - 4M_{n}^2}{s}},
\]

where \( S \) is a statistical factor to memorize the number of identical final states, \( s, t(u) = -(s/2)(1 - M_{n}^2/s)(1 \mp \cos \theta) \) are the Mandelstam variables, \( \theta \) is the angle between two incident particles in the source comoving frame. We evaluate the PDF’s \( f(x, Q) \) by CTEQ6.6m (standard MSbar scheme) [44, 45]. The calculation of total cross section \( \sigma_{\text{tot}}^{KK} \) is too cumbersome although routinely, because there are too many subprocesses and too many quark flavors. Notice that (when \( Q = 350 \text{ GeV} \)) u-quarks dominate in the large \( x \) region and gluons dominate in the small \( x \) region. So we neglect the subprocess for initial states \( s, c, b \) and anti-quarks completely in the context. We also neglect the contribution of d-quarks because their PDFs are always smaller than u-quarks; however, this reduction may not be suitable for arbitrary \( Q \). Hence, we will only calculate subprocess including u-quarks and gluons. In our calculation, \( M_{SM, n} = M_{n}^{KK} \) is assumed for simplicity, for the reason that both zero-mode on-shell mass and radiative corrections cannot alter the masses of magnitude.

The cross section is calculated in Fig. 1. We integrate out the \( \int d\theta \) to neglect the transverse momentum distribution, which is irrelevant to astrophysical scenarios. We use only the lowest order (LO) expression for quantum chromodynamics (QCD) running coupling constant \( \alpha_{S}(Q) \) in the amplitude-squared, which is sufficient for high energy regions \( Q \gg M_{Z} \). We also assume \( Q = M_{n} \) for simplicity. We do not show the uncertainties of the PDF sets in that figure; however, some tentative calculations show that error bars can be neglectable for a large range of \( Q \), and it will never affect our rough astrophysical estimations. As seen from the figure, the contribution of \( n = 1 \) state is always more important than higher excited states in the calculable regions, thus we neglect the others. The subprocess \( g + g \rightarrow g_{n}^{*} + g_{n}^{*} \) dominate the large \( \sqrt{s} \) region, and \( q + q \rightarrow q_{n}^{*} + q_{n}^{*} \) dominate.

\[ ^{1}\text{However, this assumption may need careful considerations when } \sqrt{s} \text{ is much larger than } M_{n}. \text{ We assume it because a refined calculation is far beyond our scope in this paper.} \]
the small $\sqrt{s}$ region. In magnitude, our results are consistent with \cite{13} in the regions not adjacent to $(2M_n)_+$. However, for energy just above the threshold $\sqrt{s} = (2M_n)_+$, our $\sigma$s are largely suppressed rather than mildly decreased in theirs. In addition, it seems that their $\sigma^\text{KK}(g + g \to \cdots)$ increases more softly than ours. These may only be caused by some numerical details, so we will not discuss these issues deeply. Hence $p + p$ collisions can produce the KK protons with the KK quarks or gluons in it (we will call them by a joint name $p_{\text{KK}}$). The reason why the subprocess $g + g \to g_n^* + g_n^*$ is also suitable for our purpose to produce the KK excited protons is that in the high energy collisions gluon bremsstrahlung is unsuppressed, thus the two gluons will produce two jet events to form two protons. Because we have already assumed that the KK number conservation is exact in UEDs scenario, a KK excited gluon cannot take off its KK number and cannot leave the proton for color confinement. Thus both the KK excited bosons and fermions can form the KK excited protons.

![Figure 1: Cross section for the case $n = 1, 2, 3$ and subprocess $g + g \to g_n^* + g_n^*$, $u + g \to u_n^* + g_n^*$, $u + u \to u_n^* + u_n^*$, $u + u' \to u_n^* + u_n'^*$ and $u + u \to u_n^* + u_n^*$ (Eq. (15), (17), (18), (22) and (24) in Ref. 13 respectively), as marked in the figure. The total cross sections for different $n$ are showed by thick black lines. We have assumed $q = u$ because u-quarks dominate the quark PDFs in a large range of energy scale $Q$. The other process $g + g \to q_n^* + \bar{q}_n^*$ may be important; however, we neglect it because of a lack of consistency amplitude-squared.](image.png)

The calculated cross section result has a lot of cumbersome difficulties, mainly in the high $\sqrt{s}$ regions (it has to apply) up to at least $10^{18}$ eV. The difficulties are listed as follows: (i) Because the PDFs are only applicable to $\rho_n = 4M_n^2/s > x_{\text{min}} = 10^{-8}$ in CTEQ6.6 series,
we can only calculate cross sections below energy $\sqrt{s_{\text{max}}} = 7 \times 10^{15} \text{eV}$ for the lowest KK state $M_1^{\text{KK}} = 350 \text{GeV}$, which is insufficient for astrophysical events that happen in Fermi accelerated environments. (ii) Multi-KK processes (just like the multi-$\pi$ processes in soft hadron physics) have to be considered when $\sqrt{s}$ is much larger than $M_n$. (iii) The Lagrangian we use to calculate the cross sections may only be an effective one (thus non-renormalizable in the framework of quantum field theory), hence cannot be applied to $\sqrt{s}$ larger than some cutoff energy scale (which may be not very larger than $M_1$). Problem (i) and (ii) are the main topics of §2.2. Problem (iii) is much more troublesome, thus we can only solve it completely after the emergence of the final theory (which should explain what is the physical essence of normalization). We assume their capabilities simply because (i) some smoothness considerations, and (ii) we do not have other ways to go. Other authors used the same strategies as ours, just like what in Ref. [46]. Moreover, because of the re-acceleration of our KK particles, the energetic KK ones need not to be produced by high $\sqrt{s}$ $p + p$ collisions; hence, the high $\sqrt{s}$ behaviors of our cross sections may in fact not very crucial. It should be that case at least for the naïve estimations we choose in this paper, because we can hardly distinguish the difference between the KN and DL parametrizations in Fig. 4, however, we do not know the applicability of this supposition. We leave the more quantitative discussions in the follow-up publications.

### 2.2 High Energy Behavior

As we have already mentioned before in §2.1, (i) the maximum energy a numerical total cross section $\sigma_{\text{tot}}^{\text{KK}}(\sqrt{s_{\text{max}}})$ can achieve, depends on the minimum $x$ PDFs $f(x, Q)$ can give. For CTEQ6.6 series, the believable maximum energy for the KK first-exist state $M_1^{\text{KK}}$ is only $7 \times 10^{15} \text{eV}$, which is insufficient for astrophysical purpose. Moreover, (ii) the contributions of multi-KK producing processes lack of estimations in our cross section calculations. Naïvely, one may want to continue $\sigma_{\text{tot}}^{\text{KK}}(\sqrt{s})$ by fitting any functions putting by hand, with our low energy numerical results calculated above in Fig. 1. However, if we assume that $\sigma_{\text{tot}}^{\text{KK}}$ is dominated by some special subprocess (as overpowered $g + g \rightarrow g^* + g^*$ shows in Fig. 1 in our cases), it may globally be inconsistent with the (axiomatic) Froissart-Martin bound $^3 \sigma_{\text{tot}} \leq C \ln^2 s$ for sufficiently large $s$ $^4$. Of course, everything will be fine, assuming that Froissart-Martin bound is far from being saturated.

The high energy behavior of $p + p$ total cross section is questioned in length $^4$. In this context, continuation of $\sigma_{\text{tot}}^{\text{KK}}(\sqrt{s})$ to higher energy will be done by the Kang-Nicolescu

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2We cannot saturate the bound because of our limited computation power. However, increasing the calculation load considerably can have little improvement on the result of $E_{\text{max}}$, thus we do some qualitative analyses instead in §2.2.

3Of course, the Froissart-Martin bound need not to be obeyed for a KK mass tower. Because its assumptions include that all masses are equal to the unit of mass when the energy variables go to infinity, it cannot have new states coming out when $\sqrt{s}$ increases. Spin structure is not a serious problem. Although Froissart only derived that bound for scalar particles $^4$, we can imagine it is also applicable for $2 \rightarrow 2$ quarks/gluons. Yet we can not derive it out, just for a lack of some systematical ways parallel to Mandelstam representation.
(KN) parametrization \[49\]

\[
\sigma_{\text{tot}}^{(2-2)} = A + B \log s + C \log^2 s
\]  \hspace{1cm} (2)

and the Donnachie-Landshoff (DL) parametrization (Regge fits) \[50\]

\[
\sigma_{\text{tot}}^{(2-2)} = D s^\epsilon + E s^{-\delta}.
\]  \hspace{1cm} (3)

We choose them instead of arbitrary parametrizations simply because (i) some unitarity constraints or physical backgrounds have already been considered in their derivations, which may make them more reasonable. The other reason is that (ii) the experimental result of \(p + p\) cross section for all subprocesses \(\sigma_{\text{tot}}^{p+p}\) estimated from proton-air inelastic cross section \(\sigma_{\text{tot}}^{p+\text{air}}\) \[51, 52, 53, 54, 55, 56, 57, 58\] up to \(\sqrt{s} \leq 10^{14}\) eV in cascade processes, has also been widely fitted by such kind of parametrizations \[4\]. Because \(\sigma_{\text{tot}}^{KK}/\sigma_{\text{tot}}^{p+p}\) is an important parameter in latter calculations, it will be convenient for us to compute them from the same kind of parametrizations.

The continuations of dominate subprocess \(g+g \rightarrow g_1^*+g_1^*\) by the KN and DL parametrizations, are shown in Fig. 2 and 3. However, some qualitative analyses show that, usually Eq. (3) can only construct mild concave functions in log-log diagram, just as \(\sigma_{\text{tot}}^{p+p}\) in TeV regions. The reason is that when both \(D\) and \(E > 0\), the two terms in Eq. (3) are two asymptotes with a plus sign in between. Eq. (3) can also construct convex functions in a log-log diagram when \(D \cdot E < 0\). That is, when the negative term becomes dominate, the curve will drop rapidly, which is not consistent with our purpose. Hence we simply fit \(\sigma(\sqrt{s}) = D s^\epsilon\) in Fig. 3 with the physical meaning of retaining the term analogical to pomeron exchange but dropping out the term analogical to \(\rho, \omega, f\) and \(a\) exchanges. We notice that both the KN and DL parametrizations cannot do well to fit \(\sigma_{\text{tot}}^{KK}(g+g \rightarrow g_1^*+g_1^*)\), which hints us that \(p+p \rightarrow \) (the KK states) are tremendously different from \(p+p \rightarrow \) (the (SM hadrons), hence appropriate parametrizations with reasonable physical background need to be explored. However, in this context, we use only our rough continuation for estimations below.

2.3 Collision Producing

The KK bosons and fermions can be produced by \(p + p\) collisions for high energy protons accelerated in the Fermi mechanism \[59, 60\], or a special lineage of it as the diffusive shock acceleration model \[61, 62, 63, 64, 65, 66, 67\] (see Ref. \[68, 69\] for a review), with the cross section \(\sigma_{\text{tot}}^{KK}\) already been calculated in \[2.1\] and \[2.2\]. We want to know how many of them can actually be produced in such environments.
Figure 2: Fitting of dominate subprocess $g + g \to g_1^* + g_1^*$ by the KN parametrization. $\sigma_{tot}^{p+p}$ is fitted in [58] also by the same parametrization, using the data of accelerator experiments, Akeno, Fly’s Eye, Nikolaev and GSY [51, 52, 55, 56]. The data points in the top left corner are calculated from $\sigma_{tot}^{p+\text{air}}$ in air shower experiments. The thin dashing lines are cross sections for subprocesses, as in Fig. 1 and $\sigma_{tot}^{KK}$ is shown in a thick black line. Because we can not in fact do a beautiful global fitting of $\sigma_{tot}^{KK}(g + g \to g_1^* + g_1^*)$, as shown in the figure, we fit it emphasizing particularly on (i) a smoother derivative at the endpoint, or (ii) a preferable global fitting. We will always use the first strategy in the follow-up calculations in this context.

To leave the pertinent computations for specific astrophysical sources to the later publications, and to give a more universal applicable estimation, we discuss by imitating the same assumptions in deriving the Waxman-Bahcall (WB) bound as in Ref. [70, 71]. Let $n_p$ be the number density and $E_p$ the energy of protons, the assumptions are as follows: (i) A percentage $\eta$ of the protons have an injection spectrum of $dn_p/dE_p = K \cdot E_p^{-\alpha}$ isotropically distributed from their rest mass energy $m_p = 938 \text{ MeV} \sim 10^9 \text{ eV}$ to some cutoff energy $E_{p,\text{max}}$ in the source comoving frame, while a percentage of $(1 - \eta)$ of them stay at rest in that frame. In this paper, we will assume $\eta = 1$ for simplicity. The generalization of our result to $\eta \neq 1$ cases are straightforward. The power law spectrum in energy is a natural deduction of the Fermi mechanism [59] or diffusive shock model [68]. However, for diffusive shock model, the existing derivations may need the assumption of collisionless or scatter-
Figure 3: Sameline styles have been used as in Fig. 2 but otherwise the DL parametrization. Because of some technical properties of Eq. (3) which cannot make it really looking like $\sigma_{\text{KK}}(g + g \rightarrow g^*_1 + g^*_1)$ (see the context for detail), in fact we only fit the cross section by $\sigma(\sqrt{s}) = Ds^\epsilon$.

...ing only elastically with infinitely massive objects here and there, which is unfavorable for our $p + p$ producing. Hence, a discussion of energy spectrum in a shock with influential inelastic scattering may be needed for our purpose. (ii) Sources are in the critical optical depth $\tau_{pp} = 1$, and the optical depth is independent of protonic energy. When deriving the WB bound as an upper limit, optically thin sources to $p + p$ and $p + \gamma$ reactions have been assumed. Hence, the meaning of these assumptions is that the WB bound has been saturated. When estimating the upper bound of neutrinos, $p + \gamma$ interactions may be more effective than $p + p$ interactions, thus they are given more attention. Producing of the KK particles by $p + \gamma$ reactions is also interesting (see also discussions at the beginning of §2); however, it is beyond the scope of this paper, by the lack of ready-made calculations about amplitude-squared and cross section. Here, we only consider $p + p$ interacting cases. Detailed discussions about the competition of $p + p$ and $p + \gamma$ reactions to produce both neutrinos and the KK particles are given in §5; however, a wiser way is to discuss them in association with specific astrophysical scenarios [72, 73, 74, 75]. Hidden sources with optical depth $\tau_{pp} \gg 1$ [76, 77] are very tempting for our purpose, because they can produce more KK particles by collisions. However, a big problem is how they can be accelerated and get away from the source successfully. We will discuss this situation together with some details of acceleration mechanism in §3.2, comparing it with neutrino observations in §4. Discussions linking to more realistic astrophysical environments are left to later...
publications.

If we assume that normal $p + p$ interactions which dominate $\sigma_{tot}^{p+p}$ do not change the number of protons after one time of collision, the number of the KK protons in proportion to normal protons should be

$$n_{pKK}/n_p = K^2 \int_{m_p}^{E_{p,\text{max}}} E_1^{-\alpha} dE_1 \int_{m_p}^{E_{p,\text{max}}} E_2^{-\alpha} dE_2 \times \frac{2\sigma_{tot}^{\text{KK}}(\sqrt{2E_1E_2})}{\sigma_{tot}^{p+p}(\sqrt{2E_1E_2})},$$

in which we choose $\sqrt{s} = \sqrt{2E_1E_2}$ to neglect the effect of rest mass $m_p$ and simply choose $\theta = \pi/2$. The resultant $n_{pKK}/n_p$ depending on $E_{p,\text{max}}$ and $\alpha$ is shown in Fig. 4. The cutoff energy $E_{p,\text{max}}$ in the source comoving frame can be calculated by comparing the size of the shock and the duration of an acceleration cycle, or acceleration gains and synchrotron losses of energy (see Eq. 5 and 6 for details). Sources of UHECRs should have $E_{p,\text{max}} \geq 10^{18}$ eV in the shock comoving frame, because of the reason that we have already seen a couple of UHECR events with energy $> 3 \times 10^{20}$ eV [78, 79] in the observer’s frame. Even if neglecting the energy losses by the Greisen-Zatsepin-Kuzmin (GZK) effect [80, 81] when propagating, and dividing a Lorentz factor $\Gamma \sim 300$ of the source comoving frame, we still need at least $E_{p,\text{max}} = 10^{18}$ eV. We may prefer $\alpha = 3v_d/(v_u - v_d)$ for the nonrelativistic shock cases [65], $\alpha = (3\beta_u - 2\beta_u\beta_d^2 + \beta_d^3)/(\beta_u - \beta_d) - 2 \rightarrow 38/9 - 2 \simeq 2.2$ for ultra-relativistic shock limit [67], where $v_u$ ($\beta_u$) and $v_d$ ($\beta_d$) the upstream and downstream velocities ($\beta$ factors) in the shock frame, $\alpha \leq 2$ by some statistical reasons of sources [82], or $\alpha = 2$ to be consistent with the assumption of the WB bound [70, 71]. We see that the proportion $n_{pKK}/n_p$ is not sensitive to either $E_{p,\text{max}}$ (for a large enough $E_{p,\text{max}}$, which is easy to achieve for an actual astrophysical environment) or the parametrization strategies we choose, but sensitive to the spectral index $\alpha$.

3 Accelerating and Propagating of the KK Protons

For the reason that the KK protons take charges, and may have a not-very-short lifetime (see discussions in §1) before decaying to the LKP $\gamma_1$, they can also be accelerated by the same mechanism similar to normal protons. The acceleration process is never bothered by the KK cascade decay. After being accelerated, they may also propagate through the space for a considerable distance before decaying to $\gamma_1$. If the decay processes are forbidden or suppressed sufficiently by other reasons, they may propagate as far as banging into the earth. Hence in this section, we discuss accelerating and propagating properties of the KK protons, and also propagating property of decay intermediate and final states.

3.1 The Fermi Mechanism and Diffusive Shock Model

While discussing the accelerated properties and bounds of an extreme relativistic particle, some time scales are very important. We will discuss them in the framework of the orig-
The results of KN parametrization and DL parametrization are nearly overlapped.

Figure 4: The ratio \( n_{p_{KK}}/n_p \) after one time of \( p+p \) collision for each proton in a \( E^{-\alpha} \) injection spectrum, depends on \( E_{p,\text{max}} \) and \( \alpha \). The results by the KN and DL parametrization are quite similar, thus we cannot distinguish them in this figure.

 infield Fermi acceleration system [59, 60], that is, particles are randomly accelerated by the electromagnetic turbulence. However, because of the deep connection between the Fermi mechanism and diffusive shock model, some of the results are also suitable for the latter. We will calculate these scales in the source comoving frame (for the turbulence behind the shock, it is just the shock comoving frame). To transform to the observer’s frame when the source itself is relativistic, both the size of the shock \( R \) and the energy of the particle \( E \) should multiply the Lorentz factor \( \Gamma \), e.g., \( R \to R/\Gamma \) and \( E \to E/\Gamma \).

The time scales are as follows: (i) The mean escape time \( t_E = R/c \) measures the time scale to get through the accelerating region. (ii) The Fermi acceleration time \( t_A = \eta R_L/c\beta^2 \) measures how quickly the particle gains energy, where \( R_L = E/eB \) is the Larmor radius, \( B \) is the magnetic field strength, \( \beta \) is the Alfvén velocity, and \( \eta \sim 8/3 - 50/3 \) is a factor determined by the turbulent system. (iii) Synchrotron timescale \( t_{sy} = (6\pi m^4 c^3/\sigma_T m_e^2)E^{-1}B^{-2} \) measures the energy losses of synchrotron radiation, where \( \sigma_T = 8\pi r_e^2/3 \) is the Thomson cross section, \( m = m_p \) or \( m = m_{KK} \approx M_{KK}^n \) is the mass for both normal or the KK protons respectively, and \( M_{KK}^n \) is the KK excited quark mass. Successful acceleration needs both \( t_E \geq t_A \) and \( t_{sy} \geq t_A \), thus giving two limits of maximum acceleration energy

\[
E_{\text{max}}^{(1)} \simeq \frac{cBR\beta^2}{\eta}
\] (5)
and
\[ E_{\text{max}}^{(2)} \simeq \frac{m_e^2 c^2 \beta}{\eta \sigma_T B} \sqrt{\frac{6\pi e}{\eta \sigma_T B}}. \tag{6} \]

Because of the same charges but different masses normal and the KK protons take, \( E_{\text{max}}^{(1)} \) cannot distinguish them but \( E_{\text{max}}^{(2)} \) can. Hence, to find out the KK protons by the higher energy they can achieve, we need \( E_{\text{max}}^{(1)} > E_{\text{max}}^{(2)} \) or the source to be sufficiently large, i.e.,
\[ r > \frac{\eta \ m_e^2 c^2}{eB\beta \ m_e} \sqrt{\frac{6\pi e}{\eta \sigma_T B}}. \tag{7} \]

To emphasize that up to a constant, the constraint \( t_E \geq t_A \) is equivalent to the assumption \( R > R_L \), which is a universal requirement for accelerations correlated to magnetic fields.

Notice that the formation of the cosmic-ray spectrum depends and only depends on (i) the average gain in energy per acceleration event, and (ii) the acceleration events a particle may suffer. For the original Fermi mechanism \[59\], the gain of energy is \( \delta E = (v/c)^2 E \) per collision in average (hence the energy should enhance exponentially), where \( v \) is the velocity of the reflecting obstacles (e.g., the electromagnetic turbulence). Hence, the spectrum index \( \alpha = 1 + (c/v)^2 \Delta t_{\text{ela}}/\Delta t_{\text{inela}} \) depends on the duration \( \Delta t_{\text{ela}} \) between elastic scattering (with infinite massive reflecting obstacles) and the duration \( \Delta t_{\text{inela}} \) to break down the energetic particle. For the diffusive shock model in non-relativistic shock wave cases \[69\], the average momentum (hence energy for relativistic particles) gain for isotropic distributed particles is \( \delta p = 4p(v_u - v_d)/3v_p \), and the probability of escape per shock crossing is \( 4v_d/v_p \), where \( v_p \) is the velocity of accelerated particles. To generalize this argument to relativistic shock wave cases, isotropic distributions of particles are no longer applicable, thus an intuitionistic derivation is absent. However, we may assume that the qualitative phenomena are similar. Assuming that the total inelastic cross section for the KK protons is similar to that of protons, we have \( \Delta t_{\text{inela},\text{KK}} \sim \Delta t_{\text{inela},p} \), thus both (i) and (ii) are independent of particles’ properties. The only difference between protons and the KK protons is the distinctness of their initial energy (rest mass). Hence, we have reasons to believe that the energy spectrum of the KK protons should shift to higher energy than protons by an amount of \( m_{\text{KK}}/m_p \).

The resultant \( n_{\text{pKK}}/n_p \) for interzone of fixed observed \( E \) is shown in Fig. \[5\] which is independent of \( E \) for \( m_{\text{KK}} < E < E_{\text{max}} \). Notice that \( n_{\text{pKK}} \) is enhanced by a factor of \( (m_{\text{KK}}/m_p)^{\bar{\alpha}} \) by the spectrum shift, where \( \bar{\alpha} \) is the index of the spectrum. We choose \( \bar{\alpha} = 2.7 \) or 3.0 for the overall cosmic-ray spectrum observed below and above the “knee” energy \( E_{\text{knee}} \simeq 10^{15} \text{ eV} \), and \( \bar{\alpha} = \alpha \) for some special sources. For reasonable astrophysical sources, \( n_{\text{pKK}}/n_p \) may be not too small a number, to that we have opportunities to discover them by air shower identification.

### 3.2 \( p + p \) Optical Depth Reexamined

Taking no account of the details of especial astrophysical sources, the \( p + p \) optical depth \( \tau_{pp} \) (which is assumed to be 1 in \[2,3\] and subsequently quantitative estimations) can be

\[ \tau_{pp} \simeq 10^{21} \text{ cm}^2/\text{g}, \]

If the cross section for the KK protons is smaller than that of protons, we will have the more energetic KK cosmic-rays to be observed.
Figure 5: $n_{p_{KK}}/n_p$ for fixed $E$ when $\bar{\alpha} = 2.7, 3.0$ (for diffuse flux) and $\alpha_{\text{source}}$ (for special astrophysical sources). It is not too small to be identified.

discussed by focusing on special acceleration mechanisms.

In the case of the original Fermi mechanism [59, 60], $\tau_{pp} \geq 1$ can be achieved by a sufficiently long existent duration of acceleration sources. However, high energy protons can also be broken down by $p+\gamma$ interaction or other processes, so several times of accelerations may be needed.

In the case of diffusive shock model, things are a little more troublesome. As we have already indicated in §2.3, the derivation of power law spectrum itself needs the assumption of collisionless. In the mainstream tactics for particles to cross back and forth the shock front, the downstream particles are scattered by strong magnetic turbulence behind the shock, and the upstream particles are scattered by the self-excited Alfvén waves [68]. The $p+p$ collisional disintegration has to be competitive with particle missing in downstream scattering, which makes a constraint for its feasible parameter space (because we already know some restrictions from spectrum index, etc). However, we should not be too serious in that problem, because until now there is still no consistent computation in diffusive shock model to generate the kind of waves, scatterings and acceleration phenomena [67]. Scattering with analogously energetic particles (rather than with infinite massive objects) is also a good way to change direction but avoid loosing too much energy, especially in the upstream regions where self-excited Alfvén waves themselves are in trouble.
3.3 The Modified GZK Cutoff

For the case that (i) the KK protons have too long a lifetime to decay while propagating, when discussing the propagating of protons versus them, the most (or the only) important issue is the modification of the GZK cutoff [80, 81]. Other issues while propagating, such as $e^+e^-$ pair production, cannot change tremendously the property of energy spectrum. The GZK cutoff, which is the threshold energy for protons to interact with the cosmic microwave background (CMB) photons to produce pions, as

$$E_{\text{GZK}} \simeq \frac{m_\pi^2 + 2mm_\pi}{4E_{\text{CMB}}} \approx 6 \times 10^{19} \text{eV} \left(\frac{m}{m_p}\right)$$

(8)

for $m \gg m_\pi$, is enhanced for the KK protons by $m_{\text{KK}}/m_p$ times. We have assumed that the quarks and gluons in pions are all the KK zero-mode (because the threshold energy to produce the KK excited pion is another $m_{\text{KK}}/m_\pi$ times higher), and the cross section for $\gamma_{\text{CMB}} + p_{\text{KK}} \rightarrow p_{\text{KK}}(n_{\text{KK}}) + \pi^0(\pi^+)$ is still large enough to make a cutoff. The second assumption is irrelevant because the more observable KK excited cosmic-rays lend themselves to discovery; however, the number of particles is reduced tremendously for higher energies so we have little opportunities to really meet one.

For the opposite case that (ii) the KK protons have already decayed to the LKP $\gamma_1$ (for other cases, including the decay intermediate states, the discussions are analogous; thus we will not discuss this further in this paper), they may also interact with CMB photons, as $\gamma_{\text{CMB}} + \gamma_1 \rightarrow X$. However, we do not know what $X$ really is, whether the cross section is sufficiently large to make a spectral cutoff, or what threshold energy these processes correspond to. Because the LKP $\gamma_1$ is an elementary particle (rather than the KK protons which include normal quarks/gluons), its cross sections may be much different from that of normal ones. If the process $\gamma_{\text{CMB}} + \gamma_1 \rightarrow \gamma_1 + \pi^0$ is sufficiently important, the threshold has an enhancement of $m_{\text{KK}}/m_\pi$ thus is much larger than the GZK cutoff energy, similar to Eq. (8).

4 Air Shower Identification

If we can distinguish air shower events in which primary particles are protons/ions or the KK particles (the KK protons, the LKP $\gamma_1$ or other intermediate KK excited particles which bang into the earth), we are capable of giving a tighter bound on KK contamination. Otherwise, we can only constrain it when the KK particles dominate the mass composition. The mass composition of cosmic-rays is rudimentary, partly because cascade processes with the same initial condition are not identical with each other, partly because of our insufficient knowledge of parton distribution, thus different cascade models (e.g. QGSJET [84] and SIBYLL [85]) cannot make consistent predictions with each other. An incomplete set of parameters to discriminate different compositions of cosmic-rays are as follows: (i) the elongation rate/shower maxima $X_{\text{max}}$ [86], when the energy of the particle can be decided separately by fluorescence technique, (ii) the magnitude of the fluctuation in depth of maximum $X_{\text{max}}$ [87], (iii) the number of particles $n_{\text{max}}$ at $X_{\text{max}}$, (iv) the speed of rise in $n_{\text{max}}$ [88].
(v) fraction of muons in the shower events, and some geometrically-base parameters like (vi) lateral distribution functions (LDFs), (vii) the thickness of the shower disk or (viii) the shower front curvature (see Ref. \[89\] for a review). Some of the diversities rise by the simulation results of cascade models \[84, 85\]. Hence, to identify a KK excited cosmic-ray event for our purpose, the PDFs for protons take the KK charge and corresponding air shower simulations of (i) $p_{KK}$, (ii) $\gamma_1$ or (iii) other intermediate KK excited particles are needed. However, they are beyond our scope in this paper.

Naïvely, differences mainly rise for the reason that ions are made up from constituent nucleon of relatively lower energy (typically as (iii, iv) \[88\] and the Zatsepin effect \[90, 91\]) cannot be used for reference, so a careful filtration of the derived parameters from above is needed. For some parameters, their values can only be determined statistically (such as (i) in Fly’s Eye data processing \[86\]) at present, and are thus not applicable for our purpose.

For the case of (i) the KK protons, similar total cross sections $\sigma_{tot}$ for $p+p$ and $p+p_{KK}$ are expected, because $p_{KK}$ has two/three normal valence quarks (and uncountable sea quarks) sharing energy, and $\sigma_{tot}(p+p)$ is not supposed to increase tremendously for larger center-of-mass energy $\sqrt{s}$ (see the discussions for $p_{KK}$ production in high energy $p+p$ collision in \[2\] for detail). Hence, the primary KK protons should also trigger a priori cascade processes. Differences may rise from the mass hierarchy between $p$ and $p_{KK}$, or equivalently from the energy transformation between normal and the KK quarks/gluons that should absolutely happen by color confinement, thus making the differential cross section $d\sigma(p+p)/d\Omega$ and $d\sigma(p+p_{KK})/d\Omega$ dissimilar. That may make the first few collisions much different; however, the phenomenology for the secondary showers of cascade particles like $\pi^{\pm}$, $\pi^0$, $p$, $\bar{p}$, $e^{\pm}$ and $\gamma$ are the same. Some similar work in Ref. \[92, 93\], in which the cascade process of hadronized gluinos (analogous to our KK protons with the KK quarks/gluons in them) bang into the earth is simulated, indicate that the identification is possible; however, other researchers are suspicious of their result \[94\]. It will be favorable for us, if the former result is correct, and can be naturalized to our KK proton case. Observable air shower events for the KK electrons are also expected, because it has turned on quantum electrodynamics (QED) interaction, thus the cross section should not be very small. However, the origin of the KK electrons is not discussed in this context. As a result, we can only distinguish primary protons/ions or the KK charged particles for their different cascade properties, and quantitative analyses are needed for the identifiability.

For the case of identifying (ii) the LKP $\gamma_1$, things might be a little easier. We have already had trustworthy methods to identify photons from air shower data \[95, 96\], because they interact with the geomagnetic field, thus starting the cascade process much before that of protons. Hence, if $\gamma_1$ can also interact with geomagnetic field, we can identify them easily. However, the quantitative calculations need detailed properties of $\gamma_1$ interactions.

Exotic cosmic-ray events have already been regarded, for example in Ref. \[97\]. However, they may be particularly interested as the weakly interacting low energy ones, which have their first collisions inside our scientific equipments. For our purpose, it is more pertinent to find UHE exotic air-shower events. If they would happen, they may be explained by the UEDs scenario we mention in this paper (or other new physics models), and their absence can constrain the same theoretical models.
If we cannot identify the KK charged particles’ cascade events from normal protonic/ionic ones, we can only constrain it when the KK particles dominate the mass composition. A cumbersome issue is that energy measurement is also related to the assumption of primary particles’ composition. So that even if there have already been the super-energetic KK particles (with energy largely exceeding the GZK threshold) banging into the earth, we might not discover them by our energy estimation. A wise way for energy measurements is to choose some kind of calorimetric measurements (like fluorescence light emissions [98, 99]) which are relatively model independent. Here in this paper, we assume that energy spectrum can be determined without overall departure for protons/ions or the KK particles. Thus we are capable of discovering them when they dominate the mass composition, leading to some identifiable phenomena such as a lack of the GZK cutoff, etc.

As shown in Fig. 5, if assuming $\tau_{pp} = 1$ and $\alpha = 2$, the resultant KK contamination $n_{pKK}(E)/n_p(E)$ may be as large as $10^{-5}$ to $10^{-2}$. Hence if we can identify them individually from protonic/ionic ones, it is not difficult to make a discovery. However, if we can not in fact identify them from air shower data, a larger $\tau_{pp}$ may be needed.

5 $\tau_{pp}$ Constraints from High Energy Neutrino Detectors

The $p + p$ optical depth $\tau_{pp}$ is a crucial parameter for the applicability of our methodology; however, it should be discussed in relation to special astrophysical sources. Quantitative calculations are difficult because the models for UHE particle origins themselves, such as gamma-ray bursts (GRBs), active galactic nuclei (AGNs) or supernovae (SNe), are also full of dubious issues.

However, an universal estimation can be made from the observation of neutrinos. The reason is that $p + p$ collisions can produce the KK protons, but also have branching ratios to produce $\pi^\pm$, which decay to neutrinos mainly in modes $\pi^\pm \rightarrow \mu^\pm + \nu_\mu(\bar{\nu}_\mu) \rightarrow e^\pm + \nu_e(\bar{\nu}_e) + \nu_\mu + \bar{\nu}_\mu$. Thus if we have known (or have an upper limit of) neutrino flux in earth, we can use it to restrict optical depth $\tau_{pp}$ as well as the KK cosmic-ray flux. Of course, the constraint can only be used as an upper limit, because we do not know what percentage of neutrinos are in fact produced by $p + \gamma$ interactions compared to $p + p$ interactions, or whether the KK particles can be accelerated successfully, or whether they can leave the source as easily as neutrinos do.

There are already some publications of $p + p$ collisions of the Fermi accelerated protons in astrophysical environments [72, 73, 74, 75, 76, 77, 101]. However, they are not of much use for our purpose, mainly because: (i) Their calculations are always oversimplified, because $p + p$ interactions are always unimportant, compared to $p + \gamma$ interactions in specific astrophysical environments. (ii) Their discussions do not include subprocesses opposite to

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8Other processes such as $K^\pm \rightarrow e^\pm + \nu_e(\bar{\nu}_e) + \nu_\mu + \bar{\nu}_\mu$ may also be important for our purpose, because kaons loose less energy before decaying into neutrinos, by their relatively larger mass and shorter lifetime [100]. However, we will not discuss this further in this paper.
The cross sections to produce $\pi + X$. (iii) Within some of their scenarios, one of the reacting protons does not suffer Fermi acceleration. In our scenario, it is efficient (and also favored) that both protons are accelerated. (iv) They make few attempts to calculate neutrino spectrum, or still assume $E_\nu \simeq 0.25 E_\pi \simeq 0.05 E_p$ all the same. The latter simplification is reasonable in the $p + \gamma$ cases, but not legitimate for $p + p$ cases, because usually $E_p \gg m_p \gg E_\gamma$ makes $\Gamma_p \sim \Gamma_\pi \sim (E_p + E_\gamma)/\sqrt{s}$, thus energy can be shared roughly by mass proportion.

Although spectra of thermal $p + p$ interaction have already been discussed in Ref. [102]; however, a careful calculation of energy spectrum of nonthermal $p + p$ collisions (especially in the case where both protons are accelerated by the Fermi mechanism) is absent. Briefly, only when both protons have roughly the same (direction and amplitude of the) momenta, the simplification $E_\nu \simeq 0.05 E_p$ can be used; however, the spectrum of neutrinos should be suppressed by $E^{-2\alpha}$ rather than $E^{-\alpha}$ in that case. Hence, we may guess the energy spectrum of $p + p$ neutrinos is steeper than $E^{-\alpha}$ in higher energies. Some numerical results for specific astrophysical environments show that the energy spectrum looks like $E^{-(\alpha+1)}$ [73]. We will give a more detailed universal calculation in §5.2.

### 5.1 $p + p \rightarrow \pi + X$ Cross Sections

The cross sections to produce $\pi^\pm$ in $p + p \rightarrow \pi + X$ collisions have already been calculated in [103, 104, 105, 106] (see also [107, 108] for recent developments). These parameterizations are mainly based on relatively low energy terrestrial experiments, and focus on Lorentz-invariant differential cross section (LIDCS) $E(d^3\sigma/d^3p)$ rather than total cross section $\sigma_{\text{tot}}$. Simple estimations show that they may behave well near the center-of-mass threshold energy $\sqrt{s} \geq 2m_p + m_\pi$ to produce pions, but increase too quickly to be consistent with the total cross section for all $p + p$ subprocesses $\sigma^{pp}_{\text{tot}}$ [49, 50, 58] we use in this context. Here, to give a consistent but not too cumbersome estimation from an astrophysical (rather than collider physical) viewpoint, we use Badhwar parameterization [105] when $\sigma^{\pi^+ + \pi^- + \pi^0} \leq \sigma^{pp}_{\text{tot}} \sim 40$ mb (thus $\sqrt{s} \leq 2.98$ GeV), and assume all $\sigma^{pp}_{\text{tot}}$ produce pions for larger $\sqrt{s}$ with total cross section $\sigma^{pp}_{\text{tot}}$ for a fixed proportion as the same as when $\sigma^{\pi^+ + \pi^- + \pi^0} = \sigma^{pp}_{\text{tot}}$.

The result is shown in Fig. 6. In the calculation, we assume (i) each $p + p$ interaction produces single pion, and (ii) $m_X \simeq 2m_p$ [109] when evaluating the $p_{\|,\text{max}}$ in Feynman scaling variable $x_F = p_{\|}/p_{\|,\text{max}}$. The applicability of both assumptions must be considered carefully for our purpose.

For a more accurate estimation of $p + p$ neutrino spectrum, the LIDCS $E(d^3\sigma/d^3p)$ (which depend on $\sqrt{s}$, the pion energy $E_\pi^*$ and the scattering angle $\theta^*$ in the center-of-mass frame) should be used directly to evaluate $\Gamma_\pi = \Gamma_{cm}\Gamma_\pi^*(1 + \mathbf{v}_{cm} \cdot \mathbf{v}_\pi^*/c^2)$. However, it pours oil on the flames in the further complicated calculations. Hence, we integrate out $E_\pi^*$ and $\theta^*$ by

$$
\sigma^\pi(\sqrt{s}) = 2\pi \int_0^\pi \sin \theta^* d\theta^* \int_0^\infty dp_{\pi^*}^2 \frac{p_{\pi^*}^2}{\sqrt{p_{\pi^*}^2 + m_\pi^2}} \cdot E_\pi^* \left( \frac{d^3\sigma^\pi}{d^3p_{\pi^*}} \right),
$$

In this paper, we use superscript $*$ to denote the center of mass frame, and normal characters for the source comoving frame.
and memorize \( p^*_\pi \) thereby \( \Gamma^*_\pi = E^*_\pi / m_\pi = \sqrt{p^*_\pi^2 + m_\pi^2} / m_\pi \) only by a weighting average

\[
\langle p^*_\pi (\sqrt{s}) \rangle = \frac{\int p^*_\pi (\ldots) \, d\sqrt{s}}{\int (\ldots) \, d\sqrt{s}} \tag{10}
\]

where \( \int (\ldots) = \sigma^\pi(\sqrt{s}) \). In addition, we neglect the effect of \( v_{cm} \cdot v^*_\pi / c^2 \). The result is in the top left corner of Fig. 6. We see that for \( \sqrt{s} \) not too close to \( 2m_p + m_\pi \), it is always true that \( \sqrt{s} - 2m_p - m_\pi \gg \langle p^*_\pi \rangle \). However, the result is not suitable for larger center-of-mass energy, and we lack a reasonable extension for \( \sqrt{s} > 2.98 \) GeV.

Figure 6: Numerical result of total cross section \( \sigma^\pi(\sqrt{s}) \) within Badhwar parameterization [105]. The dark region shows that the parameterization is only suitable for \( \sqrt{s} \leq 2.98 \) GeV, thus we force-fix the proportion \( \sigma^{\pi^+} : \sigma^{\pi^-} : \sigma^{\pi^0} \) for large \( \sqrt{s} \). The weighting average \( \langle p^*_\pi (\sqrt{s}) \rangle \) is in the top left corner of this figure, with the same abscissa as \( \sigma^\pi(\sqrt{s}) \).

5.2 Neutrino Spectrum

Hence, if we know the spectrum and optical depth of protons (within some specific astrophysical scenarios), the cross sections for \( p + p \) to produce \( \pi^\pm \to e^\pm + \nu_e (\bar{\nu}_e) + \nu_\mu + \bar{\nu}_\mu \), we can calculate the spectrum of neutrinos produced by \( p + p \) collisions. However, the calculation is not easily done. Even if decoupling the relationship between vectors \( v_{cm} \) and \( v^*_\pi \), and using only a weighting average momentum \( \langle p^*_\pi \rangle \) for some specific \( \sqrt{s} \), the pion
distribution function should look as complicated as

\[ n_\pi(\Gamma_\pi) \sim K^2 \int_{m_p}^{E_{p,\text{max}}} E_1^{-\alpha} dE_1 \int_{m_p}^{E_{p,\text{max}}} E_2^{-\alpha} dE_2 \times \frac{1}{2} \int_0^\pi \sin \theta d\theta \frac{\sigma^{\pi^+}(\sqrt{s}) + \sigma^{\pi^-}(\sqrt{s})}{\sigma_{\text{tot}}^{\pi^+}(\sqrt{s})} \times \delta \left( \Gamma_\pi - \frac{\sqrt{(p_\pi^+)^2(\sqrt{s}) + m_\pi^2}}{m_\pi} \frac{E_1 + E_2}{\sqrt{s}} \right) \]

(11)

where \( \sqrt{s}(E_1, E_2, \theta) \) is the center-of-mass energy in the source comoving frame. To deal with the Delta function in an easier way, we fixed \( \langle p_\pi^+ \rangle(\sqrt{s}) \equiv 0.2 \) GeV and integrate out \( \theta \) visually. However, our simplifications may underestimate the harder part of the spectrum. To estimate the absolute value of flux to normalize our calculation, we use the same assumption as in deriving the WB bound [70, 71], that is, the cosmic-rays’ energy production rate in the nearby universe is \( (E_{\text{CR}}^2 dE_{\text{CR}}/dE_{\text{CR}})_{z=0} = 10^{44} \text{erg Mpc}^{-3} \text{yr}^{-1} \), or equally \( E_{\text{CR}}^2 dE_{\text{CR}}/dE_{\text{CR}} = 5.05 \times 10^{-8} \text{GeV cm}^{-2} \text{s}^{-1} \text{sr}^{-1} \) (or 3 times larger if thinking over redshift evolution).

The resultant spectra are in Fig. 7. Our simplified estimation of \( \pi \) and \( \nu_\mu \) spectra are consistent with the numerical result \( E^{-\alpha+1} \) in Ref. [73]. In our estimation, we have already considered the effect of neutrino oscillation (detectors are only sensitive \( \nu_\mu \) flux). The WD bound can only be treated up to a constant, because if using the same treatment as ours, they should be much lower. The KK flux is estimated in a \( \alpha = 2, \bar{\alpha} = 3 \) environment (thus we use \( n_{\pi^\pm}(E)/n_p(E) \simeq 10^{-2} \)), other choices of test parameters can be read out directly from Fig. 5. We see that \( \tau_{pp} \) is not under strong constraints by neutrino observations. However, when we have more advanced neutrino detectors, they may give stronger constraints on \( \tau_{pp} \) hence the KK flux.

5.3 Other Possible Constraints of \( \tau_{pp} \) besides \( \nu_\mu \) Observations

Another possible constraint of \( \tau_{pp} \) besides neutrino flux observation is GeV – TeV gamma-ray flux observations. The reason is that \( p + p \) collisions can not only produce \( \pi^\pm \to e^\pm + \nu_e(\bar{\nu}_e) + \nu_\mu + \bar{\nu}_\mu \), but it can also produce \( \pi^0 \to 2\gamma \). Hence, the EGRET, GLAST or Swift telescope can give it a constraint. Ref. [40], in which one energetic and one rested proton collide to produce a not-very-short-lived strongly interacting massive particles, considered both neutrino and gamma-ray flux constraints. We lack of discuss the latter one, because of the fact that both (i) to avoid this paper to be too long (the discussion is very similar), and (ii) to be different from energetic neutrinos which can only be produced by violent or high \( \Gamma_{\text{cm}} \) collisions, energetic gamma-rays can also be produced by the more mild synchrotron self Compton (SSC) processes [41]. Hence the gamma-ray constraint is not as na"ive to deal with as the neutrino ones, so we leave it to later publications.
Figure 7: $\nu_\mu$ flux from $p+p \to \pi + X$ collisions. The shadow regions denote cosmic-ray flux, with the KK ones for $\tau_{pp} = 1$ and $\tau_{pp} = 10^4$ respectively. The detector information is cited from [110]. Current neutrino detectors can only give constraints that $\tau_{pp} < 10^4$.

6 Discussion and Outlook

Phenomenologies link ambitious theoretical physical models to reality, thus make physics a science [10]. Terrestrial experiments are one of the ordinary methods to constraint new physics models; however, they have limited power because of our finite energy sources on earth. Cosmology can open an extraordinary window for new physics studies; however, Big Bang (and the extreme physical environment it has) happened only once in our universe, thus makes re-enactment impossible. Astrophysical constraints always have larger scope than terrestrial experiments (because they do not have the upper threshold of maximum achievable energy); however, meaningful scenarios presently known are always only available for weakly interacting light particles. The motivation of this paper is to search for another way to construct new physics beyond the SM.

In this paper, we construct an astrophysical scenario to (dis)confirm new physics for heavy particles beyond TeV energy scale. In our scenario, the KK protons are produced by $p+p$ collisions in Fermi accelerated environments, and they themselves are accelerated by the same environments. Because they may change the compositions and properties of cosmic-ray events, air shower experiments can give a constraint to their properties.

To know whether our scenario can give meaningful constraints to UEDs (and maybe other new physics models in later researches), we make some quantitative estimations. We first investigate whether enough KK excited states can be produced by $p+p$ collisions.
We calculate the overall KK cross sections by precalculated Feynman rules and amplitude-squared, with also the CTEQ6.6m proton PDFs. Subprocesses \( g + g \rightarrow g^*_n + g^*_n \) and \( q + q \rightarrow q^*_n + q^*_n \) may be most important in quark level. Because of color confinement, the KK excited quarks and gluons should form the KK excited protons. We then calculate what percentage of the KK protons can be produced by an isotropic and power law distributed proton spectrum. For the spectral index \( \alpha = 2.0 \), \( n_{p_{\text{KK}}} / n_p \sim 10^{-10} \) as an overall contamination after one time of \( p + p \) collision. However, when considering that the KK protons are also accelerated by the Fermi mechanism, the scene is much different. It is reasonable to believe that the energy spectrum of the KK protons should shift to higher energy than protons by an amount of \( m_{\text{KK}} / m_p \), hence for some fixed energy \( E \), the KK states should contaminate \( 10^{-5} \) (for some special astrophysical sources with the spectral index \( \alpha = 2.0 \)) to \( 10^{-2} \) (for diffuse flux) of cosmic-ray events. So, if we can identify them from other cosmic-ray particles from air shower data, our method is capable of giving meaningful constraints. We notice that the GZK cutoff energy also shifts by a factor of \( m_{\text{KK}} / m_p \) to higher energy, if it still exists. Hence, observations of cosmic-ray particles with energy much above the GZK cutoff, are given a reasonable explanation by the KK particles; however, the possibilities of this kind of observation is really small, even if the GZK cutoff does not exist. We also investigate the possibilities to identify the KK cosmic-ray events by air shower data. The investigation is still superficial, because quantitative simulations are needed; however, the LKP \( \gamma_1 \) may be easy to identify, because they may interact with geomagnetic fields just like normal photons, thus make the air shower tomography much different from that of protons. Finally, we calculate the possible constraints of the KK cosmic-ray flux from neutrino detectors; however, the constraints are very loose for current scientific equipments to affect our former estimations.

It is appropriate to regard our calculations (outlined in this paper) as an “existence proof” for this kind of methodology. In fact, any charged particles beyond the SM, which are neither too light nor having a too short lifetime to suffer a Fermi acceleration, are suitable for our scenario. One immediate example is \( W^\pm_1 \) in UEDs, which we do not discuss in this paper because the extension is really straightforward. \( W^\pm_1 \) can cascade decay to \( L_1 \) or \( \nu_1 \) [14]; however, because of their analogous masses, the lifetimes of \( W^\pm_1 \) should be much longer than \( W^\pm \) in the Glashow-Weinberg-Salam theory of weak interactions. In fact, \( W^\pm_1 \) is also an intermediate state of our \( g_1 \) decay in this paper. Some lineage scenarios of the ADD model, which allow bulk bosons rather than bulk fermions [6, 7, 9], can also excite Kaluza-Klein \( W^\pm \) which suffer our accelerating. However, a careful calculation of production rates and lifetimes is absent. There are also a lot of lineage scenarios of the RS model which allow the bulk SM fields [11] [12] [13]. These models are more reasonable, because the RS model has an inherent orbifold configuration (to obtain chiral fermions), and the bulk SM fields can help us to understand some stiff physical problems like fermion mass hierarchy. Charged sparticles in supersymmetric (SUSY) models are also

\( ^{11} \)Of course, this scenario is also suitable for lighter particles. However, we can restrict the parameter space (of the endlessly emerging new models) tighter by other astrophysical/terrestrial methods. Hence, this scenario may specialize in new physics particles beyond TeV energy scale, especially just above the energy scale the best colliders can in touch.
good candidates for these kind of scenarios \[114\]. In order to explain the non-baryonic dark matter, the lightest sypersymmetric particle (LSP) is preferred to be electrically neutral; however, it is not supposed to do so. Even if the LSP is really neutral, they can also be accelerated by the Fermi mechanism if the lifetimes of the charged ones decaying to them are not very short. Whereas different from the case of the LKP \(\gamma_1\), if the LSP is gravitino or the lightest neutralino, they may hardly cause air shower processes because of their relatively small cross sections, even if they bang into the earth with tremendous energy. Gluino (which may exist as the form of gluino-containing hadron, compare to our KK proton) in split supersymmetry is also a good idea. There has already been one paper in Ref. \[40\], in which the gluinos are produced by astrophysical \(p+p\) collisions; however, (the astrophysical aspect of) our scenario has a lot of advantages than theirs, include: (i) One of the protons in their scenario stays at rest, hence the center-of-mass energy \(\sqrt{s} = \sqrt{m_p E_p}\) of \(p+p\) collisions is at least \(10^{14-15}\) eV, which is no more than \(p+\text{(air)}\) (for UHECRs to collide with the atmosphere hadrons) center-of-mass energy here in earth; however, because both of the protons in our scenario are Fermi accelerated, the center-of-mass energy of our \(p+p\) collisions should be at least \(10^{18}\) eV (we have already derived an overall Lorentz factor \(\Gamma \sim 300\), which is impossibility for any other scenarios to achieve near earth. Despite the fact that we do not know whether our ideas of (renormalizable) quantum field theory or new physics nowadays are suitable for such a huge center-of-mass energy \(\sqrt{s}\), we know that something should happen there. (ii) The maximum energy a gluino-containing hadron can achieve in their scenario is only \(10^{13.6}\) eV \[92\]; however, the maximum energy of our exotic cosmic-ray particles, can be even larger than the GZK cutoff (see \[3.3\] for a detailed discussion). Hence, because of the fact that the cosmic-ray spectrum itself has a large negative power law index of about \(-\bar{\alpha} \sim -2.7\) to \(-3.0\) (below or above the “knee” energy), for gluino-containing hadrons of energy \(E_{exo}\) produced by protons with energy \(E_p\), their content of cosmic-rays with definite \(E\) has an additional inhibitory factor of \((E_{exo}/E_p)^{-\alpha}\), which has an order of magnitude of \((10^{13.6}/E_{knee})^{2.7} \times (E_{knee}/E_{GZK})^{3} \sim 7.7 \times 10^{-19}\). If most of the UHE protons did not produce gluino-containing hadrons, the inhibition should be even stronger. Thus even if the gluino-containing hadrons are recorded by our scientific equipment (e.g., the cosmic-ray observatories), they are very difficult to be found out by such a lot of events with similar energy. However, because our charged exotic particles (KK protons in the context) are also accelerated by Fermi acceleration, their content in the UHE region should be as large as \(10^{-5}\) to \(10^{-2}\), hence not very hard to be identified. (iii) We do not really need the charged exotic particles to be longeval enough to suffer the travel from the source to earth; it is enough that their longevities are long enough to suffer an astrophysical Fermi acceleration. Hence, we do not have to worry about some adolescent (comparison with the “baby universe” era) cosmological bounds, such as the predictions of Big-bang nucleosynthesis, or cosmic microwave background. Another very interesting particle candidate for our scenario is the charged massive particles (CHAMPs) \[115\]. It is an ambitious dark matter constituent as yet (hence it is longeval to suffer a Fermi acceleration). The association of our astrophysical scenario drawing in this context and the CHAMPs, is a very interest issue; however, we leave it to the later publications.
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