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Ultra-high energy probes of classicalization

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Abstract. Classicalizing theories are characterized by a rapid growth of the scattering cross section. This growth converts these sort of theories in interesting probes for ultra-high energy experiments even at relatively low luminosity, such as cosmic rays or Plasma Wakefield accelerators. The microscopic reason behind this growth is the production of N -particle states, classicalons, that represent self-sustained lumps of soft Bosons. For spin-2 theories this is the quantum portrait of what in the classical limit are known as black holes. We emphasize the importance of this quantum picture which liberates us from the artifacts of the classical geometric limit and allows to scan a much wider landscape of experimentally-interesting quantum theories. We identify a phenomenologically-viable class of spin-2 theories for which the growth of classicalon production cross section can be as efficient as to compete with QCD cross section already at 100TeV energy, signaling production of quantum black holes with graviton occupation number $N \sim 10^4$.

Keywords: modified gravity, cosmic ray theory, cosmic ray experiments, quantum black holes

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1 Classicalization versus wilsonian physics

One of the fundamental goals of high-energy physics is to understand the nature of UV-completion of the Standard Model. In the standard (Wilsonian) paradigm of UV-completion, the new high-energy physics comes in form of weakly-coupled quantum particles that become relevant degrees of freedom at scales shorter than the weak interaction length, approximately $10^{-16} - 10^{-17}$ cm. For example, a low scale supersymmetry is a typical representative of such a Wilsonian UV-completion. As it is well-known, the above energy frontier is currently being probed by the LHC experiments.

It is natural to ask, what are the prospects of experimentally testing ideas about UV-completion in the collisions at even higher energies? One rare opportunity is provided by ultra-high energy cosmic rays (see [1] for a review). Interestingly, there are also proposals of earth-based accelerators of ultra-high energies, such as Plasma Wakefield accelerators (see [2]).

In both of these cases the potential difficulty is a relatively low luminosity, which makes the test of Wilsonian UV-completions exceedingly difficult. For example, the high energy cosmic ray collisions in the atmosphere are dominated by soft exchanges and the statistics for the events with high momentum-transfer is low. The purpose of the present paper is to identify a class of theories, based on the concept of non-Wilsonian self-completion [3] by classicalization [4], for which the ultra-high energy experiments can represent efficient probes even at relatively low luminosity, due to an efficiently growing cross-section at high energies. A potential importance of this general property for compensating the expected low luminosity of Plasma Wakefield accelerators was noted in [2].

In order to probe the short-distance physics, one needs events with high momentum-transfer. In weakly-coupled (Wilsonian) UV-completions of the Standard Model such events are predicted to be extremely rare as compared to low momentum-transfer soft events mediated by the strong interaction force. The reason is that in any weakly-coupled UV-completion of the standard model the (all-inclusive) cross section diminishes with the increase of the center of mass energy E . As a result, the cross section is dominated by soft QCD events. This is why, at least for the current statistics, the high energy cosmic rays do not represent good probes of Wilsonian UV-completions of the standard model.

Recently, a concept of non-Wilsonian self-UV-completion was introduced [3, 4], in which no new weakly-coupled physics is required above certain cutoff energy scale M_* with the corresponding cutoff length $L_* \equiv \hbar/M_*$. In these theories low energy degrees of freedom (e.g., gravitons) naively become strongly interacting at distances shorter than L_* where perturbative expansion in EL_* breaks down. However, in contrast to the Wilsonian picture, this breakdown does not imply the need for any new weakly-coupled physics that must be integrated-in at distances less than L_* . Instead, the theory cures itself in the following way. The would-be strongly-coupled particles get replaced by collective weakly-coupled degrees of freedom with an effective interaction strength suppressed by powers of $1/(L_*E)$. These collective degrees of freedom represent many-particle states of large wavelength. By large wave-length we mean the wavelengths that exceed L_* . This characteristic length is set by an energy-dependent scale which we denote by $r_*(E)$. The necessary property is the increase of $r_*(E)$ with E , so that $r_*(E) \gg L_*$ for $E \gg M_*$. As a result, the UV-theory when described in terms of collective degrees of freedom is weaker and weaker coupled and probes larger and larger distances with increasing E . This phenomenon describes the essence of what was termed as the non-Wilsonian self-completion by *classicalization* [4]. Various aspects of this idea were studied in [5]–[13]. Intrinsic feature of the classicalization phenomenon, which makes this picture phenomenologically distinct from Wilsonian completion is the efficient growth of the cross section at trans-cutoff energies as some positive power of EL_* ,

$$\sigma \propto (EL_*)^\alpha, \quad \alpha > 0 \quad (1.1)$$

This growth can be understood as the result of creation of states composed of many soft quanta, which behave more and more classically at high-energies. These are so-called *classicalons*. Let us briefly describe their essence. In the classical limit ($\hbar = 0$) classicalons represent static (usually singular) solutions of the classical equations of motion of characteristic radius $r_*(E)$ and energy (mass) E . These parameters appear as integration constants that can take arbitrary values. A well-known example of such solutions is the celebrated Schwarzschild black hole in classical general relativity. However, this geometric picture is only valid in an idealized classical limit. In reality nature is quantum and \hbar is non-zero. A quantum theory of black holes and other classicalons was developed in [9, 14]. According to the latter theory, these objects represent self-sustained bound-states of many bosons of characteristic wave-length $\lambda = r_*(E)$ and occupation number,

$$N(E) = Er_*(E)/\hbar. \quad (1.2)$$

Due to their large wave-length and derivative coupling, these bosons interact extremely weakly, with the effective coupling constant

$$\alpha_{\text{eff}} = 1/N(E). \quad (1.3)$$

Thus [14], physics of classicalons in general, and black holes in particular, is a weakly-coupled large- N physics in t’Hooft’s sense [15]. This property emerges as the result of maximal packing. The classicalons represent maximally packed states per given wave-length. The maximized occupation number density results into the oversimplification of the system and effectively converts it into a system with a single characteristics, N . In this way, classicalization replaces a would-be strongly coupled physics of few hard quanta at energy $E \gg M_*$ by an extremely weakly-coupled physics in which the same energy is distributed among many soft quanta of wavelength $r_* \gg L_*$.

This picture defines the quantum N -portrait of black holes and other classicalons. The reason for the efficient production rate of these objects in high energy particle collisions is the exponential degeneracy of micro-states that over-compensates the usual exponential suppression of many-particle states. At high energies the cross-section grows as

$$\sigma \sim (\hbar N(E)/E)^2. \quad (1.4)$$

This cross section for large $N(E)$ can be interpreted as a geometric cross section

$$\sigma \sim r_*(E)^2. \quad (1.5)$$

The relations (1.2), (1.3), (1.4) and (1.5) describe the essence of classicalizing theories. Unlike in ordinary Wilsonian case, the high energy behavior of these theories, instead of probing short distances, in reality probes large distance physics, due to the fact that the high energy scattering is dominated by production of states with large occupation number of very soft quanta.

Thus, deep-UV quantum behavior of classicalizing theories can be understood in terms of the classical IR dynamics of the same theory! For example, the behavior of deep-UV cross section can be derived by finding out the E -dependence of the $r_*(E)$ radius of a static source of mass E . The radius r_* can be defined as the shortest distance for which the linearized approximation is valid. This property simplifies enormously the predictive power of the theory for high-energies, since the dependence of σ on center of mass energy E can be read-off by solving the linearized classical equations of motion for a source of the same energy E .

However, we need to be extremely careful not to be misled by this simplification. In order to understand properly the classicalon dynamics we must continuously monitor the information obtained in an idealized classical limit ($\hbar = 0$) by translating it into the language of the underlying quantum portrait. Without this guideline, the (semi) classical picture alone can lead us to wrong conclusions. This becomes obvious, once we identify the correct classical and semi-classical limits. These limits correspond to taking

$$E \rightarrow \infty, \quad L_* \rightarrow 0 \quad r_* = \text{fixed}. \quad (1.6)$$

In addition, we may take $\hbar \neq 0$ or $\hbar = 0$ depending whether we want to be in semi-classical or classical treatment. For example, most of (if not all) the previous semi-classical analysis of the black hole physics is performed in this limit.

The quantum N portrait shows that none of these limits are correct approximations. This becomes very clear by realizing that in both limits, irrespective whether we keep \hbar finite or zero, the occupation number of quanta becomes infinite, $N \rightarrow \infty$. This immediately tells us that all the subtleties of $1/N$ expansion become hidden. The typical example of the invalidity of this approximation is the application of the semi-classical limit for micro-black holes that can be produced at LHC. It is obvious that in reality these black holes correspond to the quantum states with $N \sim 1$. Thus, to apply to their properties the semi-classical limit ($N = \infty$) gives invalid predictions, such as thermality and democracy of their decay products. In reality, the micro-black hole if accessible at LHC will behave simply as unstable quantum particles.

Likewise, for classicalons (including black holes) with large N — that as we shall argue can be observed in high energy cosmic ray experiments — the classical limit can serve as an useful guideline. But at the same time the quantum N portrait makes an essential difference

at all the scales, since it allows us to bypass huge technical complications of extracting information from the classical analysis.

For example, in the classical limit the question of black hole formation is complicated with all possible technical subtleties, such as the question of horizon formation. Since in such a limit the quantum nature is completely hidden, the only remaining characteristics are geometric and one has to be extremely precise to prove the black hole formation. These technicalities, if the only guideline used is the geometric picture, make the outcome of high-energy scattering inconclusive.

Instead, the large N portrait sheds a very powerful light on this process. It makes clear that all the geometric issues, such as horizon-formation are secondary if we are interested in the question of formation of large- N quantum states. For large N , this question can be answered reliably in $1/N$ expansion, which is an excellent approximation for the cases of interest.

With the large- N quantum portrait in mind, the questions whether the large- N states are formed can be reliably answered already by linearized classical analysis properly translated into the underlying quantum language.

Equipped with the above knowledge, in the present note we shall focus on a class of theoretically-motivated classicalizing theories that can be of interest for high energy cosmic rays as well as for other high-energy but low luminosity experimental searches, with the signatures that can be cross correlated with the signatures at LHC.

The class of theories of our interest represents UV-modification of Einstein's gravity that classicalize at energies above $M_* \sim \text{TeV}$ and thus are theoretically motivated for solving the hierarchy problem. They can be viewed as a generalization of large extra dimensional scenario [18], but we shall abandon the geometrical framework allowing ourselves to characterize the dynamics in terms of degrees of freedom and their occupation number. As we shall see, this approach since it liberates us from the frame of geometric thinking, allows us to uncover a much wider class of phenomenologically-interesting possibilities.

We shall show that in the most efficiently classicalizing scenario, the high-energy growth of the cross section can be as steep as

$$\sigma(E) \simeq \pi (E^2 L_*^4). \quad (1.7)$$

This approximately-quadratic growth takes place until the energy $E_c \equiv (L_*^2 m)^{-1}$, where, m is the mass of the new spin-2 degrees of freedom. Above E_c the power-law growth of the cross section changes to a logarithmic growth according to

$$\sigma(E) = \pi m^{-2} \ln^2 \left(\frac{EL_*^2}{\sqrt{\sigma(E)/\pi}} \right). \quad (1.8)$$

We estimate the phenomenological lower bound on the scale m , and show that for $M_* \sim \text{TeV}$ it can be as low as $\sim 20\text{MeV}$. This lower bound comes from the combined constraints of supernova cooling and the big bang nucleosynthesis.

To conclude, we show that in the interval of phenomenologically interesting energies, classicalons with occupation number as large as $N \sim 10^4$ can be produced in high energy cosmic ray collisions with atmosphere at energies as low as 100 TeV, with their production cross-section competing or even exceeding the typical low-momentum-transfer QCD cross-section $\sigma_{QCD} \sim 100\text{mb} \sim 250\text{GeV}^{-2}$.

Of course, the possibility of micro black hole production in above-TeV-energy collisions is not new and was already suggested in [19] and in a number of subsequent papers (see,

e.g., [20]). This suggestion was based on combining the intuition about the classical black hole formation in ultra-Planckian energies pioneered in earlier papers [21–23] with the new input of possibility of lowering the fundamental Planck mass down to TeV energies [18]. Therefore, most of the previous studies were done within the classical geometric frameworks such as [18], which cannot capture the underlying quantum large- N structure of black holes. By abandoning the pure geometric limit we manage to identify classes of theories with more efficient growth of the cross section. The purpose of the present paper is not to confront these models with any particular experimental data, but to point out their phenomenological viability and their potential experimental importance for ultra-high energy experiments, both cosmic ray as well as the earth based accelerators.

A very interesting representative of such an earth-based ultra-high-energy accelerator is a currently-planned proton-driven wakefield accelerator [2]. These experiments have a potential of going significantly above the current accelerator capacities, but with relatively low luminosity. However, for probing classicalizing theories the low luminosity is not a fatal problem since the large cross section can make up for it. Hence, such ultra-high energy experiments represent potentially very interesting probes for the classicalization phenomenon.

2 Spin-2 classicalizers

We shall discuss the phenomenological implications of classicalizing theories in which the classicalizer fields are spin-2 bosons. Our aim is to identify the most efficient low-energy classicalizers that could be of potential interest for cosmic ray physics. These theories are in the same universality class as the original large extra dimensions in which fundamental scale of gravity is $M_* \sim \text{TeV}$ [18]. In the latter class of theories, above the energy scale set by the inverse compactification radius, the theory is effectively described by Einstein gravity in $4 + d$ dimensions, with the fundamental Planck length being L_* . Such a theory classicalizes above the energy scale M_* . According to the large- N quantum portrait, the classicalons of this theory are many-graviton states of wavelength,

$$\lambda = r_*(E) = L_* (E L_*)^{\frac{1}{1+d}}, \quad (2.1)$$

and occupation number

$$N = \hbar^{-1} E r_*(E) = \hbar^{-1} (E L_*)^{\frac{2+d}{1+d}}. \quad (2.2)$$

Putting it shortly, black holes are graviton Bose-condensates. Given the wave-length, the quantum interaction strength between graviton pairs is given by

$$\alpha_{gr} = \hbar (E L_*)^{-\frac{2+d}{1+d}}. \quad (2.3)$$

Notice that all these quantities can be expressed in terms of N as,

$$\lambda = N^{\frac{1}{2+d}} L_* \quad \alpha_{gr} = 1/N. \quad (2.4)$$

In the limit $\hbar = 0$, $L_* \rightarrow 0$, $N \rightarrow \infty$, these objects can be viewed as $4 + d$ -dimensional Schwarzschild black holes of radius r_* . Thus, physics of black holes is the same large- N physics that was used by t'Hooft for $SU(N)$ gauge theories, except the role of N is played by graviton occupation number.

However, the large- N quantum portrait of black holes makes it clear that overemphasizing the geometric meaning is not always very useful. First, for moderate values of N the

geometric picture is simply not a good description. Secondly, it illustrates that restricting the analysis to the configurations that have well-known geometric interpretation in the $\hbar = 0$ limit, leaves aside a huge class of potentially interesting cases which can be easier to access in high energy experiments and could have more spectacular experimental signatures than the "canonical" cases which have been studied in the geometric classical limit.

Therefore, the correct attitude in this search would be to adopt the fully quantum mechanical view in which the main role is played by the degrees of freedom and their occupation number, with a geometric interpretation being secondary.

Gravity viewed as a quantum field theory is a theory of a dynamical metric, $g_{\mu\nu}(x)$, which can be represented as a background metric $\langle g_{\mu\nu} \rangle$ plus a small perturbation, $\delta g_{\mu\nu}$,

$$g_{\mu\nu} = \langle g_{\mu\nu} \rangle + \delta g_{\mu\nu}. \quad (2.5)$$

We shall be interested in dynamics of the sources on asymptotically-Minkowski spaces, which is an excellent approximation for the problems we shall study.

In such a case, the metric created by a static source sufficiently far from it, is approximately Minkowskian and can be described as,

$$g_{\mu\nu} = \eta_{\mu\nu} + \delta g_{\mu\nu}, \quad (2.6)$$

where $\delta g_{\mu\nu} \ll 1$. The key approach of the quantum N -portrait is to think of such a geometry as a quantum state (a sort of Bose-condensate) that is characterized by the occupation number N , the wavelengths of gravitons λ and by their interactions strength α_{gr} . Classicalons are then the special states in which all these parameters assume values that are fully determined by N , in such a way that a self-sustainability condition is satisfied. For example, the quantum counterparts of Schwarzschild black holes are given by (2.4).

The dictionary between the quantum and classical characteristics can be established as follows.

Metric perturbation $\delta g_{\mu\nu}$ can be written in terms of linear perturbations of canonically-normalized fields that can be classified by their transformation properties with respect to the Poincare group. Approaching the source from infinity, we shall measure $\delta g_{\mu\nu}$ which can be found by solving the appropriate linear equations, in which it is sourced by an effective energy-momentum tensor $T_{\mu\nu}$ of energy E . Assume that this approximation breaks down at some distance $r \equiv r_*$ for which $\delta g_{\mu\nu}$ becomes order one. We shall call the corresponding radius, r_* the *classicalization* scale.

The quantum mechanical meaning of r_* is that we are dealing with a quantum-mechanical state with graviton occupation number given by (1.2) and the characteristic wavelength,

$$\lambda = r_* \quad (2.7)$$

The term *classicalization radius* has an obvious meaning. It is enough to notice that $N \gg 1$ whenever r_* exceeds the Compton wavelength of the source (\hbar/E). That is, we are dealing with a large occupation number state that can approximately be treated as a classical object. In a quantum Universe like ours, this is the only consistent meaning of classicality.

Also, r_* acquires a classical geometric meaning in the formal limit $\hbar \rightarrow 0$ and $N \rightarrow \infty$. For example, in Einstein theory, in such a limit r_* becomes a geometric Schwarzschild radius. However, we shall not use this meaning. Treating r_* quantum-mechanically shall allow us to bypass unnecessary details emerging in exact classical limit, which in any case represent an

idealized approximation of the real world and are only approximately applicable to a realistic scattering processes.

Following the large- N quantum portrait of black holes and classicalons, for us r_* will be the scale (the characteristic wave-length) for which the interaction of N -gravitons becomes strong and the system enters the regime in which a self-sustained bound-state starts to form.

This picture allows us to draw important conclusions about the quantum dynamics of this self-sustained bound-state without going beyond the linearized approximation on the IR side.

The key is to identify the energy dependence of the cross section. As we shall show this is determined by the number and the couplings of gravitational species that $\delta g_{\mu\nu}$ propagates. The number of new graviton species n which is a fixed input characteristics of a given theory, should not be confused with the occupation number, N , of graviton quanta in a given black hole. The latter number depending on the black hole mass and takes different values for different black holes in one and the same theory.

The most general expansion of the *linearized* metric perturbation in terms of graviton species has the following form,

$$\delta g_{\mu\nu} = \frac{1}{M_P} \sum_m g_m^{(2)} h_{\mu\nu}^{(m)} + \frac{1}{M_P} \sum_m g_m^{(0)} \eta_{\mu\nu} \phi^{(m)}, \quad (2.8)$$

where $h_{\mu\nu}^{(m)}$ and $\phi^{(m)}$ stand for spin-2 and spin-0 degrees of freedom respectively. The constants $g_m^{(2)}$ and $g_m^{(0)}$ parameterize the strength of their coupling relative to the zero-mode Einstein graviton ($h_{\mu\nu}^{(0)}$), for which we have $g_0^{(2)} = 1$.¹

Consider now a localized static source of some characteristic wave-length $\lambda = R$. Classically, such a source can be approximated by a static energy momentum tensor smeared over a spherical shell of radius R , $T_{\mu\nu} = \delta_\mu^0 \delta_\nu^0 \frac{E}{4\pi R^2} \delta(r - R)$. For such a source, the linear graviton perturbations can be easily obtained by solving appropriate linearized equations which are well-known and we shall not display. In particular, for the Newtonian components we get,

$$h_{00}^{(0)} = \frac{L_P E}{2r}, \quad h_{00}^{(m)} = \frac{2}{3} g_m^{(2)} \frac{L_P E}{r} e^{-mr}, \quad \phi^{(m)} = g_m^{(0)} \frac{L_P E}{r} e^{-mr}. \quad (2.9)$$

Substituting this expression into (2.8) we get for the total metric perturbation the following expression,

$$\delta g_{00} = \frac{L_P^2 E}{r} \sum_m \rho_m e^{-mr}, \quad (2.10)$$

where ρ_m parameterize the relative strengths of massive spin-2 states with respect to the massless graviton. That is, $\rho_m \equiv \frac{2}{3} (g_m^{(2)})^2 + (g_m^{(0)})^2$, whereas $\rho_0 = 1/2$.

It is clear that the sum in (2.10) measures the relative strength of gravity with respect to pure Einstein gravity, and can be written as,

$$\frac{\delta g_{00}}{\delta g_{00}^{(Einstein)}} = \sum_m \rho_m e^{-mr}. \quad (2.11)$$

¹Notice, that since we couple the metric perturbation only to the conserved energy momentum tensor, the would-be contribution from the derivatively-coupled scalars of the form $\partial_\mu \partial_\nu \phi^{(m)}$ as well as the contributions from spin-1 states, vanishes. These are therefore neglected in the expansion (2.8).

The quantum-mechanical picture described by the metric perturbation (2.10) is that a source of energy E induces a state of graviton occupation number,

$$N = \sum_m N_m = (E^2 L_P^2) \sum_m \rho_m e^{-mR}. \quad (2.12)$$

Thus, each graviton of mass m and coupling ρ_m contributes the partial number $N_m = N_0 \rho_m e^{-mR}$, where $N_0 \equiv E^2 L_P^2$ is the occupation number of the massless gravitons.

In our parameterization the classicalization self-completion threshold is set by the source of the Compton wave-length L_* that creates an order-one metric perturbation at the distance $r \sim L_*$. That is, $\delta g_{00}(r = L_*) = 1$. Or equivalently, in the quantum language this is a source for which the occupation number of gravitons, N , exceeds one. In other words, such a source marks the boundary between the one-particle and multi-particle states.

Let us see, what relation this imposes on the spectrum of gravitational messenger species. Since each graviton species of mass m contributes only at distances $r < 1/m$, at any given distance r we are allowed to only take into the account modes that are not heavier than $1/r$. Thus, for the estimate we can truncate the sum at $m = 1/r$. At the same time we can approximate the exponential factors for all the light modes by one. The sum then simplifies to,

$$\delta g_{00}(r) = \frac{r g}{r} \sum_{m=0}^{m=1/r} \rho_m. \quad (2.13)$$

Now, the criterion that sources heavier than L_*^{-1} must classicalize reads

$$\delta g_{00}(r = L_*) = \frac{L_P^2}{L_*^2} \sum_{m=0}^{m=1/L_*} \rho_m = 1. \quad (2.14)$$

The bottom-line is that the cumulative strength of all the new degrees of freedom at the scale L_* must be L_*^2/L_P^2 times the strength of Einstein's gravity. This condition only constraints the overall sum of coupling strengths as given by (2.14). This constraint can be accommodated in different ways. For example, one can introduce many weakly-coupled gravitons or few strongly coupled ones. By changing coupling strengths ρ_m and the number of gravitons we can scan an entire landscape of classicalizing spin-2 theories.

Sources of energy $E \gg M_*$ classicalize and represent N -particle states. This applies both to static sources as well as to scattering two-particle states with center of mass energy E . Our task is to single out the class of theories for which the classicalization radius can exhibit the most efficient growth at phenomenologically-interesting low energies i.e above $M_* \sim \text{TeV}$. This task is simple to accomplish. Each member of the graviton tower of mass m_j and of strength ρ_{m_j} contributes into the classicalization process at distances shorter than its Compton wavelength m_j^{-1} and decouples at larger distances. Since the integrated over-all strength of gravitons is fixed by (2.14) the most efficient growth of $r_*(E)$ with E will be achieved if we make all the new gravitons maximally light.

The classicalization radius of any source of energy $E \gg M_*$ is determined by the relation,

$$\delta g_{00}(r_*) = \frac{E}{M_P^2 r_*} \left[\sum_j e^{-m_j r_*} \rho_{m_j} + 1 \right] \simeq 1 \quad (2.15)$$

We now wish to find such consistent values of m_j and ρ_j that would ensure the largest possible value of r_* for a given mass E . This is obviously achieved in the case when most of the graviton masses are pushed down as much as possible. At any given energy E the gravitons contributing efficiently into the growth of r_* will be the ones that satisfy,

$$m_j < m = \frac{M_*^2}{E}. \quad (2.16)$$

The extreme case is achieved when all the gravitons open up at the same universal mass m . Since what is important is the integrated strength, we can take all the graviton species to be coupled with the uniform strength set by their inverse number, n ,

$$\rho_m = L_*^2/L_P^2 \frac{1}{n}. \quad (2.17)$$

In other words, for a given center of mass energy E , the largest classicalization radius $r_*(E)$ is achieved when all the n graviton species open up at the common Compton wavelength $m^{-1} \gg L_*$. The $r_*(E)$ radius of such a classicalon is then determined by the condition:

$$\frac{E}{M_*^2} \frac{1}{r_*} \left[e^{-mr_*} + \frac{M_*^2}{M_P^2} \right] = 1 \quad (2.18)$$

In order to translate this result into the relation between the classicalon production cross-section (σ) and the center of mass energy (E), all we need to do is to identify $\sigma = \pi r_*^2$ and plug it in (2.18). This identification gives us the relation,

$$E = M_*^2 \frac{\sqrt{\frac{\sigma}{\pi}}}{e^{-\sqrt{\frac{\sigma}{\pi}}m} + M_*^2/M_P^2}. \quad (2.19)$$

We can further simplify the above relation by noticing that for the values of interest $M_* \sim \text{TeV}$, the term M_*^2/M_P^2 is of order 10^{-32} and is absolutely negligible. The value of $r_*(E)$ is then determined by the condition

$$\frac{E}{M_*^2} \frac{e^{-mr_*}}{r_*} = 1 \quad (2.20)$$

or equivalently,

$$E = M_*^2 \frac{\sqrt{\frac{\sigma}{\pi}}}{e^{-\sqrt{\frac{\sigma}{\pi}}m}}. \quad (2.21)$$

Which in terms of cross-section gives us (1.8). For $m < (EL_*^2)^{-1}$ the value of r_* is given by

$$r_* \simeq EL_*^2 (1 - mEL_*^2), \quad (2.22)$$

which in terms of cross-section gives us leading expansion of (1.8) in mEL_*^2 . The latter expression is no longer applicable after the center of mass energy reaches $E_c = (mL_*^2)^{-1}$, or equivalently the r_* radius reaches m^{-1} . Above this energy we have to use the expression (2.21) (which implies (1.8)), which indicates that the growth of cross section becomes logarithmic,

$$\sigma = \pi m^{-2} \ln^2(mEL_*^2). \quad (2.23)$$

As it is clear from (2.19), this logarithmic growth persists till the energies of order $E \sim M_P^2/m$, after which one enters into the regime in which classicalization is taken over by

the Einsteinian massless graviton via production of Einsteinian black holes. Obviously, this regime is way beyond the reach of current cosmic ray experiments. So the signature that is of interest for the current observations is given by the cross section (1.8). As we shall show in the next section, the phenomenological lower bound on m is as low as $\sim \text{MeV}$, allowing the classicalization cross section to compete with or even dominate the QCD cross section at the energies around $\sim 100\text{TeV}$.

3 Phenomenological constraints on m

As we have shown above, the efficiency of the cross section growth in deep-UV is determined by the Compton wave-length of the gravitons. This is a very peculiar characteristics of classicalizing theories and is not present in weakly-coupled Wilsonian UV completions. For example, the behavior of the cross section at energy $E \gg L_*^{-1}$ instead of being determined by the strength of the interaction at distance \hbar/E , is rather determined by the strength at distance $r_* \gg L_*$. The gravitons of Compton wavelength shorter than r_* effectively do not participate in formation of classicalons. Thus for the efficient growth of the cross section at energy E it is important to have as many light gravitons as possible with the Compton wavelength exceeding $r_*(E)$. However, graviton masses and couplings are constrained from below by various observations. In case of large extra dimensions this constraints have been studied in great detail [24]. We shall now apply the similar analysis in our more general case.

The number and couplings of gravitons available at each mass level m are constrained by various high-energy processes as well as by the tabletop measurements. We shall now briefly evaluate these constraints.

3.1 Star cooling

We shall consider the star-cooling constraints first. These constraints come from the fact that gravitons can be produced in the interiors of the stars and contribute to their cooling. The requirement that this is not disturbing the usual cooling rate, places the constraint on the masses and couplings of the graviton species.

We can consider two different regimes depending whether mean free path of produced gravitons is longer or shorter than the star radius R_{star} . In the first case the gravitons escape freely, whereas in the second case they can deposit the energy back into the star before escaping. We shall consider a free-escape regime first.

3.1.1 Free escape regime

In such a regime the star-cooling rate is essentially given by the graviton production rate, since probability of the re-capture is small.

For a star with an interior temperature T , the cooling rate due to emission of massive gravitons is,

$$\Gamma_{star \rightarrow graviton} \sim \frac{T^3}{M_P^2} \sum_j \rho_j e^{-m_j/T}, \quad (3.1)$$

(where for simplicity we do not make distinction between the spin-2 and spin-0 cases). The exponent comes from the Boltzmann suppression. This rate must be less than the standard cooling rate. The most stringent bound comes from supernova 1987A, with the core temperature $T = 30\text{MeV}$. In order to make a quick estimate, we can use the knowledge from the

previous analysis (such as production of axion [25] or Kaluza-Klein gravitons [24]) indicating that the cooling rate into new states should not exceed the quantity

$$\Gamma_{\max} \sim \frac{T^3}{10^{18} \text{GeV}^2}. \quad (3.2)$$

Demanding that (3.1) does not exceed (3.2) gives the following constraint,

$$\sum_j \rho_j e^{-m_j/30 \text{MeV}} < 10^{20}. \quad (3.3)$$

Let us recall, that according to (2.11) the sum entering on the right hand side of (3.3) measures the strength of gravity relative to pure Einstein at distances $r = (30 \text{MeV})^{-1}$. This means that the star-cooling constraints could still tolerate modified gravity that at distances $r \sim (30 \text{MeV})^{-1}$ could become 10^{20} -times stronger relative to standard Einstein gravity. With this constraint in mind let us evaluate the case which could give the largest possible cross section in the range of energies that can be potentially probed by high energy cosmic ray experiments.

The two choices of the parameters that tend to enhance the classicalon production cross section are: 1) Pushing the scale M_* as low as possible; and 2) increasing the Compton wavelengths of gravitons as much as possible. So the maximal effect is achieved when almost all the new graviton species open up at the same mass scale m .

In such a case, the ratio of the supernova cooling rate into the gravitons to the maximal allowed cooling rate (3.2) becomes,

$$\frac{\Gamma_{\text{star} \rightarrow \text{graviton}}}{\Gamma_{\text{standard}}} \sim \left(\frac{10^9 \text{GeV}}{M_*} \right)^2 e^{-m/30 \text{MeV}}. \quad (3.4)$$

LHC experiments already set the lowest bound on scale M_* in the $\sim \text{TeV}$ range. Substituting this lower limit in (3.4) and demanding that the ratio must be below one we arrive to the lower bound on m of few hundred MeV.

3.1.2 Energy recapture regime

We shall now consider a different regime in which the individual gravitons are strongly enough coupled so that their mean free path is less than the star radius. Such a situation, for example, can take place when gravitons that are produced in a hot inner core of the star decay into the ordinary particles before reaching the surface. In such a case they deposit energy in the interior of the star before escaping and this suppresses their contribution into the cooling. Let us evaluate constraints for such a regime. The lifetime of a graviton of mass m_j and coupling ρ_j/M_P^2 produced at temperature $T \gtrsim m$ in the interior of the star is

$$\tau \sim \frac{T}{m_j} \frac{M_P^2}{\rho_j m_j^3}. \quad (3.5)$$

The coefficient T/m_j comes from a relativistic gamma factor. This lifetime sets the distance that gravitons travel before decaying. Demanding that this distance is less than the radius of a proto-neutron star, $\tau < R_{\text{star}} \sim 10 \text{km}$, we get the following bound on the mass,

$$m_j^4 \gtrsim \frac{T M_P^2}{R_{\text{star}} \rho_j}. \quad (3.6)$$

For the case, in which all the graviton species have the same mass m , we can translate the above bound in terms of their number, n , using the relation (2.17),

$$m^4 \gtrsim \frac{TM_*^2 n}{R_{star}}. \quad (3.7)$$

This has a clear physical meaning. For less number of graviton species, in order to accommodate the same collective strength, the individual couplings must be stronger. Thus, the probability of re-capture for each graviton becomes higher. As an extreme case we can take $n = 1$. In this case taking $M_* \sim \text{TeV}$, we get the following re-capture bound,

$$m \gtrsim \text{MeV}. \quad (3.8)$$

Notice that in approximating the mean free path by (3.5) we assumed that the particles into which gravitons decay are off-shell (due to thermal bath) by a sufficiently small amount so that the decay is still allowed. This is a reasonable approximation for the current estimate.

To summarize, the supernova gives the following constraints on the scale m in two different regimes. In a free-streaming regime, when gravitons are many and weakly coupled, the lower bound is about of few 100MeV. In the recapture regime, when gravitons are less and strongly coupled, the bound is $\sim \text{MeV}$.

3.2 Other astro-cosmo constraints

With the above bounds being satisfied the other astro-cosmo constraints (e.g., such as nucleosynthesis or the diffused gamma ray background) are either automatically satisfied or are much milder.

Let us consider for example a big bang nucleosynthesis (BBN) bound. This bound can be estimated from the requirement that graviton production rate at the BBN temperature $T_{BBN} \sim \text{MeV}$, should be subdominant to the expansion rate of the Universe which in that epoch is $H_{BBN} \sim T_{BBN}^2/M_P$. This requirement gives the bound,

$$\sum_j \rho_{m_j} e^{-m_j/\text{MeV}} < 10^{22}. \quad (3.9)$$

This bound is milder than (3.3) and thus is automatically satisfied when supernova cooling constraint is met in the free-streaming regime. For example, when all the graviton species have a common mass $\sim 100 \text{ MeV}$.

In the recapture case, when the number of graviton species is smaller and they are stronger coupled, BBN bound is slightly stronger than the supernova bound and gives $m \gtrsim 20 \text{ MeV}$.

Given the constraints from supernova and BBN the constraints from the diffused photon background are easy to accommodate. This constraint usually comes from a possible late decay of the gravitons [18].

Even if produced with a maximal abundance, for the value of m even at its lowest bound 20 MeV, the lifetime of gravitons can be easily made shorter than the age of the Universe. Indeed, for number of graviton species being n , the lifetime of graviton of mass m is,

$$\tau_{gravitons} \sim M_*^2 n / m^3. \quad (3.10)$$

Requiring that this is less than the age of the Universe, $\tau < H_0^{-1} \sim 10^{28} \text{ cm}$, we get a very mild upper constraint on the total number of gravitational species $n \lesssim 10^{32}$, which in any case is satisfied due to the constraints on the number of gravitational species implicit in equation (2.14).

3.3 tabletop measurements

Other constraints on the Compton wavelengths of new graviton species come from tabletop measurement of deviations from Newton's law. These experiments have been probing distances down to the micron range and are placing bounds on gravity competing forces of various strengths.

In case when all the new gravitons are heavier than the star cooling constraint, the Compton wavelengths are too short to be of interest for the above-mentioned table top measurements.

However, as we have shown above, the combined star-cooling and BBN constraints allow existence of longer Compton wave-length gravitons as long as their cumulative strength is less than $\sim 10^{20}$ times the strength of Einstein gravity. This bound leaves a huge room for the existence of lighter gravitons that could be of interest for table top experiments. The absolute upper bound on Compton wave-length of any new graviton interacting with approximately the same (or larger) strength as a massless graviton, comes from the torsion balance experiment and is 0.5 mm [26]. Shorter scale experiments are sensitive to stronger forces.

4 Concluding remarks

The point of this note was to outline the potential importance of classicalizing theories for ultra-high energy experiments with relatively low luminosity, such as the high-energy cosmic ray or plasma Wakefield accelerators [2]. We have identified some phenomenologically-consistent theories that exhibit a most efficient growth of the cross section at the energies of experimental interest.

We have stressed importance of underlying quantum picture [14] according to which the black holes (or classicalons in general) are Bose condensates of maximal (for a given size) occupation number N . All the characteristics of this system can be understood in terms of N . For the phenomenological studies in the ultra-high cosmic ray energy range, this quantum picture is important because of at least two reasons. First, it allows to map the linear analysis on an underlying quantum portrait, which would be impossible in the classical limit $N = \infty$, $r_* = \text{fixed}$. Secondly, it allows to cross correlate expected events at LHC with the ones in ultra-high energy searches. For example, in the considered classicalizing models within the energy interval between $\sim 1 - 100$ TeV we scan the classicalons states with occupation number ranging in an interval $N \sim 1 - 10^4$. Thus, we predicts, that the two-jet events at LHC that are expected to come from the decay of the lightest classicalon resonances ($N \sim 1$) must be cross-correlated with the multi-jet events in higher energy collisions.

We have identified a phenomenologically-viable class of classicalizing spin-2 theories, and have shown that the current observational constraints put a limit on the saturation point of the cross-section-growth at about $\sigma \sim (20 \text{ MeV})^2$. This cross section dominates over the soft QCD cross section and thus represents an interesting potential target for the experiments with high energy but low luminosity collisions, such that are expected in cosmic rays or at plasma wakefield accelerators.

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