Why Unparticle Models with Mass Gaps are Examples of Hidden Valleys

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Hidden valleys, hidden sectors with multi-particle dynamics and a mass gap, can produce striking and unusual final states at the LHC. Unparticle models, hidden-sectors with conformal dynamics and no (or a very small) mass gap, can result in unusual kinematic features that indirectly reflect the conformal dynamics. When sufficiently large mass gaps are added to unparticle models, they become hidden valley models. Predictions using unparticle propagators alone overlook the most striking signals, which are typically of hidden-valley type. Inclusive signatures often cannot be predicted from unparticle dimensions, and exclusive signatures are often visible and can be spectacular. Among possible signatures are: Higgs decays to pairs of particles that in turn decay to two quarks, leptons or gauge bosons, possibly with displaced vertices; Higgs, top, and neutralino decays to more than six particles; resonances below an “unparticle” continuum which produce multi-body final states; etc. The Stephanov model is deconstructed, reconstructed, and shown to be a hidden valley model. Some effects of strong dynamics on hidden valley observables, not predictable using unparticle methods, are discussed, including resonances, reduced flavor symmetry breaking, reduced supersymmetry breaking, and a strongly enhanced hidden parton shower.

1. INTRODUCTION TO HIDDEN VALLEYS AND UNPARTICLES

The “hidden valley” scenario \[1\] envisions a hidden sector coupling weakly (but not too weakly) to the standard model, with a multi-particle production mechanism and a mass gap (or an equivalent; see below.) It encompasses a very wide class of models, since almost any gauge theory from the last forty years, with a mass gap added, will fit within this category. A few examples of such models were given in \[1\]. The phenomenological signatures, which can include novel and sometimes spectacular signals at particle accelerators, were explored in \[1, 2, 3, 4\], but much more study is needed. The main point is that some of the particles which are produced in abundance in the hidden sector can decay back to standard model particles on detector time-scales, creating a directly observable signal at the LHC.

The “unparticle” scenario \[2\] (see also \[5, 8, 10\]) envisions a hidden sector coupling weakly (but not too weakly) to the standard model, with conformal dynamics (itself a multi-particle production mechanism) and, in the original papers, no mass gap. It encompasses a wide class of models; essentially any conformal or near-conformal field theory known from the last forty years will do. In this case the effects described in \[1, 2, 3\] are absent, because the particles in the hidden sector distribute their energy into massless modes (or, if there is a tiny mass gap, into states with lifetimes too long to observe.) Then one must rely on less spectacular but no less interesting inclusive processes, such as kinematic distributions in events with missing energy, and effects on production of small numbers of standard model particles.

Both scenarios are motivated in part by the fact that hidden sectors coupling to the standard model at or near the TeV scale are common in non-minimal beyond-the-standard-model theories, for example in supersymmetry breaking models and in many string theory constructions. The Randall-Sundrum-2 scenario \[8\] is dual (by gauge-string duality) to such a model. Other recent examples include the original Twin Higgs \[11\] and Folded Supersymmetry models \[12\]. The other motivation for considering the two scenarios is that they both give unusual signals at the LHC. In the hidden valley case, these signals often create experimental challenges that have been considered very little, if at all. To avoid missing these signals altogether in the complex environment of the LHC, it is vital that these signatures be studied.

Soon after the term “unparticles” was introduced, mass gaps were added to unparticle theories, for example in \[13, 14, 15\]. Somewhat surprisingly, the fact that unparticle models with mass gaps are identical to hidden valley models — ones that have conformal dynamics above their mass gap — seems not to have been recognized.

In particular, all known conformal field theories in four dimensions are gauge theories, which all have the parton shower dynamics that plays an important role in \[1\]. With a mass gap, many of these models also have hadronization and/or cascade decays, as again is common in hidden valleys \[1\]. The predictions of hidden valley models are qualitatively the same, and many of the experimental implications similar, independent of whether the multiparticle dynamics above the mass gap is conformal or non-conformal, and whether it is weakly-coupled or strongly-coupled. The main differences are ones of degree. Thus an unparticle model with a mass gap satisfies all the criteria for a hidden valley model, and should have the same predictions, even though the literature, up to now, has indicated otherwise.

In this paper I will show how to make the hidden valley phenomenology of these models more obvious. I will emphasize that focusing on the inclusive “unparticle” phenomenology often overlooks the main predictions for experiment. The most dramatic signals typically come from exclusive final states that cannot be studied using current unparticle methods. Even the most striking inclusive signals are often not of unparticle type. In any
case, unparticle propagators are not sufficient for prediction; much more detailed information about the conformal sector and about the breaking of conformal invariance is needed.

I will now summarize the main points of this paper. These are a combination of general cautionary remarks and predictions of LHC signatures that may arise in this context.

• Within the unparticle language for describing conformal field theory, there are two critical issues, often ignored, that strongly impact phenomenology. These are the imaginary part of the unparticle propagator, and unparticle interactions, including both self-interactions and interactions with other composite operators in the conformal field theory. With a mass gap, these effects become even more central, typically dominating the phenomenology. (Throughout this paper, I will generally refer to a mass “gap”, but in fact only a mass “ledge” — where one or more particles get stuck because they cannot decay within the hidden sector, forcing them to decay via the standard model — is necessary for many of the conclusions.)

• Once a mass gap is introduced at the scale \(M\), conformal invariance and the dimensions of operators cannot predict much of the physics in the vicinity of \(M\). The details of the hidden sector and of the precise form of conformal symmetry breaking are critical. Not even the dominant cross-sections, much less distributions of kinematic variables for the most common events, can be predicted from unparticle propagators. Unparticle methods are reliable only on tails of distributions, where all the kinematics lies high above \(M\).

• The hidden sector often becomes much more visible once a mass gap is introduced and the sector becomes a hidden valley. Hidden valley phenomenology can be spectacular. However, the range of visible phenomena that can occur in a hidden valley is enormous, as emphasized in [1,2,3], and cannot be predicted from the dimensions of the operators in the conformal field theory above the mass gap. Instead, many more details of the hidden sector are required in order to make predictions for the LHC, and in particular, to determine whether the hidden sector phenomenology is invisible, is visible but challenging, or is visible and spectacular.

• As an example, the mixing of the Higgs boson with an “unparticle” can easily lead to spectacular Higgs decays, such as Higgs decays to four leptons, Higgs decays to eight or more partons, Higgs decays to two or more particles which decay with displaced vertices, etc. This is for the same reason as in [2,16,17]; see also [18,19,20]. Again, the details cannot be predicted from the dimensions of operators; specific knowledge of the hidden sector and of its conformal symmetry breaking are required for any predictions.

• It has been proposed that a mass gap in a conformal hidden sector will lead to a tower of narrow states, with lifetimes decreasing as a power law as one goes up the tower. This is only true in an extreme situation. Instead, it is far more likely that all but one or two of the lightest states will decay to one another with short lifetimes not predictable from unparticle methods. Only the lightest one or two states in each tower will be extremely narrow and may have long enough lifetimes to give displaced vertices. This is exactly analogous to what one sees in QCD, or in pure Yang-Mills theory [21], and is also precisely what one expects in a confining hidden valley model [1].

Meanwhile, strong dynamics in a hidden sector can have other, non-unparticle, effects.

• Strong couplings can lead to sharp resonances at a scale very roughly of order \(M\), but with a distribution of masses and widths that cannot be predicted from any unparticle method.

• Large anomalous dimensions (of operators not necessarily coupled to the standard model) can suppress global symmetry breaking and/or supersymmetry breaking in the hidden sector, with potentially observable consequences. It is possible that supersymmetry might be discovered first in the hidden sector before evidence for it in the standard model sector is convincing.

• Operators of high spin and large anomalous dimension, which cannot be given low-dimension couplings to the standard model, can strongly affect the hidden parton shower, which often plays an important role in hidden valleys; see for example [1]. The effect is that high energy events, instead of producing two jets of hidden-sector-particles, produce a large number of soft hidden-sector-particles, in a distribution that probably is quasi-spherical. This has important experimental consequences.

We first review the hidden valley scenario and its predictions. We briefly review unparticles and discuss the effect of adding conformal symmetry breaking. In Sec. [11] we explore some of the physical phenomena which can arise when conformal symmetry is broken, emphasizing its diversity. We reconsider [14] and come to different conclusions. In Sec. [11] we deconstruct the Stephanov model of unparticles [15] and reconstruct it, coming to very different conclusions from [15] about the effect of conformal symmetry breaking. Finally in Sec. [11] we discuss some non-unparticle effects of strong coupling in the hidden sector, and their possible consequences for observable signatures at colliders.
FIG. 1: In the hidden valley scenario, a hidden sector couples at or near the TeV scale to the standard model sector. In the simplest hidden valleys, a barrier limiting production of hidden-sector particles will be breached in the near future. The number of particles increases through a multi-particle production process in the hidden sector. A mass gap prevents decays within the hidden sector, allowing hidden-sector particles to decay to visible particles, often with long lifetimes due to the barrier. Events with high multiplicity and/or displaced vertices naturally result.

A. The Hidden Valley Scenario

In the hidden valley scenario, a model must have three ingredients, illustrated in Fig. 1:

- a coupling through a “communicator” ("mediator", "portal") to a hidden sector,
- a multi-particle production process in the hidden sector, and
- a mass gap or “ledge” which prevents some particles in the hidden sector from easily decaying to lighter particles in the hidden sector

The specific realization in any model may take many forms.

The communicator could be any neutral particle, including

- neutral gauge bosons, such as $Z$ or $Z'$ bosons [1],
- Higgs bosons [1, 2],
- neutralinos [3],
- right-handed neutrinos.

Communication could also be generated through a loop of particles charged under both standard model and hidden sector gauge groups [1].

The multi-particle production mechanism might include [1]

- cascade decays of massive particles,
- parton showering (but note this need not be QCD-like parton showering)
- hadronization (but note this need not be QCD-like hadronization)

In QCD, let us note, all three mechanisms are operative, and all three may be active in the hidden sector. However, any one of the mechanisms is sufficient for the predictions below. Note also that in the gauge-string correspondence (often called AdS/CFT [22, 23, 24] or AdS/QCD [25, 26, 27, 28, 29, 30], with RS-type models [8, 9] arising in a certain limit), these dynamical processes can be represented through equivalent processes in five dimensions, as we will review in Sec. III.

Finally, a mass gap (or its equivalent) could arise from several sources:

- explicit masses from, for example, supersymmetry breaking or the Higgs expectation value.
- confinement (of a form similar to, or different from, QCD) [1]
- the Higgs mechanism (which is electric-magnetic dual to confinement)
- compactification of an extra dimension (which is often dual to confinement or the Higgs mechanism)

The standard model and its supersymmetric extension exhibit examples of the first three. Meanwhile the fourth, if the extra dimension has a warp factor, is known in some cases to be a crude dual description of confinement in QCD, and in some cases to be an exact dual description of confinement effects (or of the Higgs mechanism) in some non-QCD-like gauge theories.

The reason to take such a large scenario with so many classes of models in it is that it makes a general class of novel predictions. This is analogous to the way that one gathers many supersymmetric models, extra dimensional models and little-Higgs models together because of their basic mechanisms, which immediately imply missing energy signals and heavy partners for some or all standard model particles. The predictions made in [1], which follow from the general structure of hidden valley models, are

- new light neutral states, decaying to the standard model through a variety of decay modes,
- long lifetimes for the new states, including therefore the probability of substantial missing energy and the possibility of highly displaced vertices,
- abundant high-multiplicity final states in high energy processes, with large event-to-event fluctuations in multiplicity, visible energy, event shape, and other quantities.

Other possible phenomena include
• non-standard multi-body decay modes for the Higgs boson, including ones already discussed \cite{2,16,19,20} and beyond, including the possibility of discovery channels involving highly displaced vertices \cite{1,2};

• new decay modes \cite{3} for the lightest R-parity-odd (or KK- or T-parity-odd) particle in supersymmetric or extra-dimensional or little Higgs models, and indeed in any model with a new global symmetry; again these decay modes offer several possible sources of highly displaced vertices and high-multiplicity final states.

Since these signals arise so easily and generally, and yet many of them do not appear often or at all in the standard array of most-studied models — technicolor, supersymmetry and its cousins, extra dimensions or little Higgs models — they pose potentially new challenges and opportunities for the Tevatron and LHC experiments. A little exploration of the experimental literature and internal notes shows many of the associated issues have not yet been addressed. More than anything else, it is this point which makes the hidden valley scenario important to consider.

B. The Unparticle Scenario and the Effect of an Added Mass Gap

The unparticle scenario is equally general. In its original form it requires only two ingredients:

• A coupling of one or more gauge-invariant local operator $O_{sm}$ in the standard model to one or more gauge-invariant local operators $O$ in the hidden sector.

• A conformal field theory in the hidden sector with no mass gap, or a very low one.

The predictions of the scenario include missing energy signals with unusual kinematic distributions for the visible particles that depend on the dimension of the operator $O$, and new production mechanisms and interference effects for two-to-two scattering of standard model particles \cite{6,31}, as well as in multi-particle production \cite{32}.

The addition of a third ingredient — conformal symmetry breaking that generates a mass gap or its equivalent — turns these theories into hidden valley models, in particular ones with ultraviolet-conformal dynamics. The general predictions of the models are therefore, not surprisingly, the same. The detailed predictions of course depend on exactly what the conformal sector contains, and especially, as we will see, on how conformal symmetry is broken. This is not easy to discern from the literature, however.

For example, in \cite{15} a strong breaking of conformal invariance was shown to be inevitable in a large class of unparticle models, and the physical effect of a mass gap was modeled. In \cite{14} the mixing of a Higgs boson with an unparticle sector with a mass gap was explored. In \cite{15} a mass gap was added to a toy unparticle model to illustrate some theoretical points, and also the physical implications of the mass gap were briefly discussed. In all these cases, as we will see, the full story was not told, and the neglected visible signatures are of the type expected in hidden valleys, as described in \cite{1,2,3,4}.

II. HIDDEN VALLEY PHYSICS OF UNPARTICLES WITH A MASS SCALE

In this section I will consider the addition of conformal symmetry breaking at a scale $M$ to an unparticle sector, mostly concentrating on the case where $M$ is of order or greater than a few GeV. Using a few toy models, I will illustrate how the dominant phenomenology typically cannot be predicted using unparticle techniques. This is especially true at the Tevatron and LHC: the rapidly falling parton distributions bias the phenomenology toward the lowest accessible energies, where conformal symmetry breaking is most manifest. Even if the sector is largely invisible, its production and interference effects are highly variable and cannot be predicted from operator dimensions alone. Instead the details of the conformal sector and its conformal symmetry breaking allow for enormous variety. Moreover, the sector itself may become visible, with all the phenomenology of \cite{1,2,3}. If so, the resulting exclusive events are often the leading observable at the Tevatron and LHC, and often even more
so at the ILC.

I will mostly focus on a popular scenario in the unparticle literature, where a hidden sector is coupled to the standard model though the Higgs boson \[13, 14\]. Many conclusions drawn are not specific to this case, although my presentation will highlight certain details of this particular unparticle coupling. I will consider some other cases briefly in Sec. [IIC].

### A. Three Toy Models of the Hidden Sector

In order to illustrate various physical points, I will introduce three useful toy models of the hidden sector. Model A, hidden scalar QED, will have the advantage of extreme simplicity and great phenomenological riches, at the cost of being fine-tuned and requiring us to remain at weak coupling. Model B, a scalar Banks-Zaks theory, will be only slightly more complicated and will have even richer physics; it is strictly conformal and fixed points might exist at stronger coupling. However it too is fine-tuned. Nevertheless, its concepts can be extended to fermionic Banks-Zaks theories, which are not fine-tuned. Finally, model C, a supersymmetric Banks-Zaks theory, is a theory very similar to model B, as far as its unparticle physics, and suffers no fine-tuning. It has physics far too rich to fully explore in this paper, as it includes all of the phenomenological possibilities of models A and B as a small subset.

Now I will describe the models in more detail, and how their unparticles couple to the standard model. Except in section [IIC], where I will briefly consider other examples and show that the conclusions are similar, I will focus on the case where the unparticle couples to the Higgs boson. This case is already sufficient to uncover the exquisite complexity of hidden valley phenomenology.

Model A is weakly-coupled scalar QED: a theory of a photon plus \(N_f\) massless scalars \(\phi_i\) of charge 1. This theory has a small beta function if \(\alpha N_f\) is small, as I will assume, so it violates conformal invariance by a small amount. But small violations of conformal invariance have small effects on unparticle predictions, as the reader will easily verify in the discussion below. In this theory, the mass operators \(\phi_i^\dagger \phi_j\) develop small negative anomalous dimensions and can serve as unparticles of dimension just below 2.

Model B is a scalar Banks-Zaks fixed point with an \(SU(N)\) gauge theory and \(N_f \sim 22N\) scalars in the fundamental representation. Again the mass operators will serve as unparticles. Although fixed points at strong coupling may exist, I will only consider the weakly-coupled cases.

I have chosen scalars rather than fermions here because I want some weakly-coupled examples, and many authors claim unparticles only make sense for \(d_\text{O} < 2\), to avoid divergences. I disagree, but do not want to be distracted by this controversy here. The mass operator in a scalar Banks-Zaks theory has \(d_\text{O} = 2 - \text{order}(\alpha N/\pi)\), an unparticle by any measure.

The third toy model addresses the problem of naturalness which is present in models A and B. Model C is a supersymmetric \(SU(N)\) Seiberg fixed point \[\gamma\], not necessarily weakly coupled, with \(N_f\) flavors of quarks \(\psi_i, \tilde{\psi}^i\) and squarks \(\tilde{\phi}_i, \tilde{\phi}^i\), in superfields \(\Phi_i, \tilde{\Phi}^i\). When \(N_f \sim 3N\) the theory is a Banks-Zaks point and is weakly coupled. This theory is natural, and its squark-antisquark bilinear \(\tilde{\phi}_i \phi^j\) is an unparticle of dimension \(3 - 3N/N_f\), which approaches 2 from below as \(N_f \to 3N\).

As mentioned earlier, the first two theories are highly fine-tuned; they involve scalars \(\phi\) and gauge bosons, and both \(\phi^\dagger \phi\) and \((\phi^\dagger \phi)^2\) are relevant operators. However, one can check that all results obtained using these models apply also to the third case, which is a natural theory without fine-tuning. Many also apply to fermionic Banks-Zaks models, albeit for a fermion bilinear unparticle, whose dimension is close to 3 if the coupling is weak, though it may be much smaller at strong coupling.

### B. Important General Observations

Before I put these models into action, I would like to make a few observations which are very important for many applications to phenomenology, and are very typical in conformal or near-conformal sectors.

- All physically reasonable conformal field theories have non-zero three- and higher-point correlation functions.
- All conformal field theories have composite operators of higher spin and dimension; these will not couple to standard model fields in the Lagrangian but can still play an important role in the physics.
- Many conformal field theories, including these toy models for \(N_f > 1\), have flavor symmetries under which the unparticles transform as one or more multiplets.

On the first point, note the following. In models A and B the three point function for \(\langle \tilde{\Phi} \Phi \Phi \rangle\) is obviously non-zero; in perturbation theory this is a non-vanishing loop diagram. Importantly, this three-point function does not go to zero as the hidden-sector gauge coupling goes to zero. Even as the hidden gauge theory becomes free, the unparticles do not become free. Composite operators have interesting \(n\)-point functions even in free theories, and these “unparticle interactions” must not be neglected, especially once conformal symmetry is broken. We will see this shortly. The only way around this is to take a conformal gauge theory with \(N\) colors and take \(N\) strictly to infinity. At any finite \(N\), the \(1/N\) corrections change the physics drastically, and we will see later that even \(N \sim 10000\) is not large enough that one can ignore these interactions for phenomenological predictions.

On the second point, there are operators in the hidden sector which must be present. Obviously these include...
the stress tensor and any conserved flavor currents. Less obviously, any gauge theory has high-spin high-dimension operators; for example, the so-called “DGLAP” operators are always present. These cannot serve as unparticles, since the couplings to the standard model would be highly irrelevant, but they can affect the phenomenology, as we will see in Secs. III and IV.

On the third point, models A, B and C all have multiple unparticles, as do many reasonable conformal field theories, transforming under a flavor symmetry. The Higgs may couple to one linear combination of these operators. This is the case in all of our examples with \( N_f > 1 \). If there are no interactions other than the gauge interactions (and their supersymmetric partners, where appropriate), then the first two theories would have \( U(N_f) \) symmetry, and model C has \( SU(N_f) \times SU(N_f) \times U(1) \) symmetry. In models A and B, then, the operators \( O_i^j \) break up into two subclasses: the operator \( O = \sum_i O_i^j \), which is a singlet under \( U(N_f) \), and the remaining operators, which form an adjoint of \( U(N_f) \). Since the \( U(N_f) \) currents commute with the dilatation operator, the members of the adjoint all have a dimension \( d_A \), and the singlet in general has a different dimension \( d_0 \). (For example, in \( N = 4 \) Yang-Mills, theory, there are six scalars in an SO(6) global symmetry; the adjoint bilinear has dimension 2 for any coupling while the singlet has a positive anomalous dimension; see [33] for an application to the hierarchy problem.) In model C, however, the operators \( O_i^j = \phi_i \phi^j \) are in the bi-fundamental representation of \( SU(N_f) \times SU(N_f) \), and all share the same dimension, unless additional interactions are added.

The interaction with the Higgs may break these flavor symmetries, splitting the degeneracies within the multiplets and allowing processes that would be otherwise forbidden. Such effects are very important for the phenomenology, and very model-dependent, as we will see shortly. In general, if multiple standard model operators couple to multiple unparticles, they will couple to different linear combinations. One must keep track of the breaking of any global symmetries, as it will affect the observed phenomenology.

Thus an unparticle coupled to the standard model cannot be treated in isolation. It may transform non-trivially under exact or approximate flavor symmetries in the hidden sector. It is an interacting object, both with itself and with other composite operators that may not be coupled to the standard model; in fact a free unparticle of dimension above 1 is not consistent. It will interact with the energy-momentum tensor and with higher-dimension and higher-spin operators in the theory. These interactions become extremely important when conformal symmetry is broken.

C. The Predictions of Conformal Invariance

Now let us couple the Higgs boson to the operator \( O = \sum_i O_i^j \) (the flavor trace of the mass operator) in any of the toy models. In models A and B we simply add \( f H^i h O \) to the Lagrangian; see for example [13] for a discussion. Recall that the mass operator has a small negative anomalous dimension (at least if the gauge coupling is the largest coupling in the theory, which we will assume). Of course this interaction will badly destabilize the fixed point in our first two toy models, since relevant operators \( O \) and \( |O|^2 \) will be induced. But these models are indeed toys, so we accept some severe fine-tuning, as in the standard model, in return for simple-minded illustrations that can easily then be generalized to realistic cases.

In model C there are two choices. We could introduce a term supersymmetrically, by writing \( \Delta W = y H_u H_d S + \zeta SO \), where \( S \) is a singlet and \( H_u, H_d \) are the two Higgs doublets; this induces the term \( |y H_u H_d + \zeta O|^2 \) in the Lagrangian. Alternatively we could introduce \( \Delta W = y H_u H_d O \), though by unitarity this is an irrelevant operator, so its coefficient might be suppressed. In either case we will also assume that supersymmetry breaking adds additional terms to the scalar potential. Rather than treat these carefully, we will speak in more general terms about the low-energy dynamics, which will be sufficient to illustrate the complexity and diversity of possible phenomenology. Also, note that supersymmetry breaking itself can break conformal invariance in the hidden sector. In this case, a small coupling to the Higgs boson does not necessarily imply a low conformal breaking scale.

Now when the Higgs gets an expectation value, a “tadpole” term \( f v^2 O \) will appear in the Lagrangian. This is simply a mass term for the scalar fields \( \phi_i \). Obviously the conformal symmetry of the hidden sector is broken. Also, through the term \( f v h O \), the physical Higgs and the operator \( O \) will mix, which allows the process \( gg \rightarrow h \rightarrow O \), Fig. 3. This process was studied in [14].

![FIG. 3: An unparticle mixing with the Higgs boson is produced in \( gg \) collisions.](image)
\( \phi^i \) and \( \phi \) must come together to form a loop, as in Fig. 4. Alternatively, if we consider the imaginary part of the \( \mathcal{O} \) propagator, this contains the \( gg \rightarrow h^* \rightarrow \phi^i \phi^i \) process.

The changed dimension of \( \mathcal{O} \) is also clear from this viewpoint. The departing \( \phi \) and \( \phi^i \) are attracted to each other by the hidden gauge interactions, modifying the cross-section and cause it to fall faster than the \( 1/s \) behavior that would be expected if the scalars were free.

**FIG. 4:** In our toy models at weak coupling, the unparticle propagator is a \( \phi \) loop, corrected by hidden gauge boson exchange, and the imaginary part of the unparticle propagator contains the process \( gg \rightarrow \phi \phi^i \).

In models A, B and C at weak coupling, not only do virtual gauge bosons at the production vertex change the scaling law, through ultraviolet effects, but also they can be emitted in the production process, along with \( \phi_i \) and \( \phi^i \), as in Fig. 5. How should we represent this in terms of unparticles? It is another contribution to the imaginary part of the \( \mathcal{O} \) propagator, which in fact will get contributions from an infinite number of processes involving the scalars and gauge bosons. If we only consider exclusive questions, we can treat the \( \mathcal{O} \) propagator as a black box and ignore exactly what is going on inside the imaginary part. But we will see in the next section that the phenomenology depends crucially on looking inside this box.

The two-point function of \( \mathcal{O} \) must be modified at energies where conformal symmetry breaking is important, and its precise form cannot be determined without detailed understanding of the hidden sector. In [13] a form for this two-point function was proposed, valid for any \( d < 2 \). The functional form has a sharp cutoff at some value of \( q^2 \); it would be visible in experiments, as noted in [34]. But we can immediately see that this is not the form which applies for any of our toy models.

**D. The Unbroken Phase: A First Look**

There is an important question at the first stage, which is whether the breaking of conformal symmetry, which gives the \( \phi_i \) a mass, might also give them expectation values. Let us first assume that we are in an “unbroken phase” where the \( \phi_i \) are massive but the hidden gauge symmetry is not broken.

In models A and B, at weakly coupling, we can immediately see what happens. Let \( m \) be the physical mass of the \( \phi \) field, which at weak coupling differs only slightly from \( \sqrt{f} v \). The calculation of the cross-section is almost that of a free theory, which means that there is a phase-space suppression of the cross-section at \( q^2 \) just above \( m^2 \). Here the continuum production smoothly ends. But this is not all: there are \( \phi^i \phi_i \) bound states, which give a set of resonances below the cut. These are not infinitely narrow, because they can both radiate and annihilate to hidden gauge bosons, so only a finite number of resonances are resolvable. They are weak and closely spaced if the hidden gauge coupling is small, but strong and spread out for more strongly coupled theories. Thus the cross-section for “unparticle” production for \( d \mathcal{O} \) near 2 might resemble Fig. 6. Observation of these resonances allows for an alternative measure of strong coupling, one complementary to the measurement of the power law obeyed by the falling cross-section.

But if the Higgs couples with different coefficients to the various \( \phi_i \) fields — if it has small couplings to unparticles \( \mathcal{O}^i \) other than \( \mathcal{O} \equiv \sum_i \mathcal{O}_i^i \) — then the breaking of conformal invariance may be yet more complicated. Rather than one threshold with one set of bound states, there may be several. This could lead to a cross-section given by the bold curve in Fig. 7. The curve in Fig. 6 is shown as a thin curve on the same plot; note that these differ strongly near the peak cross-section, yet both match the unparticle prediction perfectly at high energy. These cross-sections are simple partonic cross-sections;
FIG. 6: One possible shape for the cross-section $\sigma(s)$ to produce hidden sector particles, versus partonic energy; parton distribution functions are not included here. Note the smooth turnoff, due to decreased phase space, as $s$ decreases, and the resonances from hidden-sector $\phi^+\phi$ bound states.

the cross-section at LHC, Fig. 8 (a log-linear plot!), where the parton distributions are folded in, strongly de-weights the high-energy region, where the unparticle prediction is valid.

FIG. 7: Another possible shape for the cross-section $\sigma(s)$ versus the partonic energy, with the curve of Fig. 6 shown for comparison; parton distribution functions are not included here. Both curves are appropriate for $N_f = 4$; the thick curve shows the result of having four different masses, one somewhat lighter than the other three, while the thin curve has equal masses. Note the curves agree exactly at high energy but differ greatly near the point of maximum cross-section.

Moreover, in this case the unparticles may not be invisible. If any of the diagonal flavor symmetries are broken, either by the Higgs couplings to the unparticles — for example, if there is an $H^1 H O_{12}$ coupling, even a very small one — or by flavor-violating $O$ or $O^2$ terms in the Lagrangian, then nothing forbids the decay $\phi_2 \rightarrow h \phi_1$. Here the Higgs may be on- or off-shell, depending on whether $m_2 - m_1$ is larger or smaller than $m_h$. It is possible that this visible decay may be overwhelmed by an invisible flavor-changing coupling of the $Y$ boson, namely $\phi_2 \rightarrow Y \phi_1$, but this is obviously model-dependent. Certainly we should not assume that unparticle production is invisible, and we must reconsider the experimental implications and whether this helps, or hinders, a measurement of the unparticle cross-section.

That the decay $\phi_2 \rightarrow h \phi_1$ is possible is clear as day in the language of the scalars and gauge bosons in models A, B and C, at least when the coupling constant is weak and the unparticle has dimension just below 2. Yet how unclear it becomes when one tries to write it in terms of the gauge invariant operators $O^j_i$! The otherwise obvious process becomes quite obscure.

FIG. 8: A log-linear plot of $\log[\sigma(\sqrt{s})]$ versus $\sqrt{s}$ at the LHC, for the same two models shown in Fig. 7. The gluon distribution functions greatly enhance the region of greatest difference between the two models.

We can try to write it down anyway, as an exercise, and also to see that the process does not vanish when the coupling becomes stronger and perturbative intuition fails. In unparticle language, there is a $\langle O_{22} O_{21} O_{12} O_{11} \rangle$ four-point function, even in the limit that the hidden gauge coupling goes to zero. This induces a process $gg \rightarrow h^* \rightarrow O_{22} \rightarrow hh O_{11}$. A cut through this process includes the process $gg \rightarrow hh \phi_1 \phi_1^\dagger$, Fig. 9. This would be very small in a conformal field theory coupled to the Higgs, but when conformal symmetry is broken, the analytic structure of the $O^j_i$ propagators is drastically altered, making the process now unsuppressed. The existence of flavor-symmetry-breaking spurions, and the change in the analytic structure of the $O^j_i$ propagators (whose cuts now end at finite timelike values of $q^2$, not at $q^2 = 0$), are enough to guarantee that processes with Higgs emission are enhanced compared to the case with unbroken conformal symmetry.

1. Possible Signatures

Despite the fact that there are visible effects, this model may pose a considerable challenge for the LHC. At best one obtains the following interesting but difficult signatures:

- If $m_h < 2 m_2$ and $m_2 - m_1 > m_h$ then in production of $O_{22} (gg \rightarrow \phi_2 \phi_2^\dagger)$ one might observe a final state
with two Higgs bosons plus MET from the invisible \( \phi_1 \) and \( \phi'_1 \). The best channel might well be \( b\bar{b}\tau^+\tau^- \) plus MET.

- If \( m_h < 2m_2 \) and \( m_2 - m_1 < m_h \), then the decay of \( \phi_2 \) must go via off-shell Higgs bosons; the rates are suppressed. It is possible the potentially invisible channel \( \phi_2 \rightarrow \phi_1 Y \) will dominate, but this may also be suppressed. The likely signal would involve soft nonresonant \( b \) quark pairs or \( \tau \) pairs plus MET.

- If \( m_h > 2m_2 \), then the decay \( h \rightarrow \phi_2 \phi_2 \) is allowed, followed by the decay \( \phi_2 \rightarrow h^* \phi_1 \) with an unknown branching fraction. From such decays, the visible energy may be quite small, so triggering on \( gg \rightarrow h \) may be impossible. In vector boson fusion events one might observe forward jets plus soft jets or leptons and MET from the two \( \phi_2 \) decays. Similar signals would arise in \( Wh \) and \( Zh \) production.

All of these modes are challenging or perhaps even impossible at the LHC. Indeed these are signals best found at an ILC, or perhaps the Tevatron. But in any case, the hidden sector is not generally invisible. In particular, a search aimed at unparticles assumed to be invisible, such as for events with a \( W \) and large MET with no central jets, might throw away the signal.

This is not the complete list of possibilities, and at larger \( N_f \) the possible signatures multiply, as cascade decays ensue, giving high-multiplicity final states. These signatures are a minimal type of hidden valley-like phenomenology, in which a \( \phi_i \) that is trapped on a “ledge” cannot decay rapidly within its own sector, and instead decays through a visible-sector Higgs boson. This kind of decay should be common to many models with multiple sterile scalars mixing with the Higgs boson [18]; however these signatures do not seem to have been explored.

Let us compare this conclusion with that of [14], in which a Higgs boson coupling to the hidden sector was considered. In figures 3 and 6 of [14], cross-sections for unparticle production take the form of a pole below a continuum, if the Higgs is lighter than \( 2m_2 \), and of a broad resonance inside a continuum, if the Higgs is heavier. In the second case we recognize this as the typical behavior of a particle mixing with a continuum – a pole mixing with a cut. (See also [18].) In [14] it was stated that the signal is invisible, unless the unparticle itself can decay by mixing back through the Higgs.

Why were the conclusions of [14] so different from those of this section? A particular model for the unparticle was employed, a specific regulated form (Eq. 2.8 of [14]) of the unparticle/Higgs coupling was introduced to assure stability of the vacuum, and unparticle interactions were neglected. These choices are simple, but they are not characteristic of typical hidden-sector gauge theories. Most importantly, their lack of realism precisely assumes away all the signals discussed above. (Moreover, the assumptions made cause the unparticle to develop an expectation value. This has other important effects that we will see in a moment, in Sec. III F.)

In more realistic models, there are likely to be interesting resonances, as in Fig. 6 below the continuum, and possibly more structure, as in Fig. 7 if global symmetries are broken. And if there is any flavor mixing, as could be induced by a mismatch between the Higgs couplings and hidden sector couplings, then we expect decays within the hidden sector via Higgs boson emission, as in Fig. 9 and in the bullet points above. The rates and branching fractions depend upon the rest of the hidden sector, not specified in [14].

In sum, the hidden sector is easily made visible once conformal symmetry breaking occurs. But we are not done, by any means, at least not in general models. In model A, we may have identified all the phenomenology of the unbroken phase, since the theory may simply consist of massive scalars and a massless hidden photon. The massless scalars have a \( \mathbb{Z}_2 \) symmetry which forbids the lightest state \( \phi_1 \) from decaying to the standard model, and the same is true of the hidden photon \( Y \), which also therefore remains invisible. But in models B and C we must address what happens at lower energy to the so-far massless hidden gluons. We will do this in Sec. III F.

Before we do this, let us consider a completely different possibility.

E. The Broken Phase

After conformal symmetry breaking, the theory may also end up in a “broken phase”, where the gauge symmetry is broken. The presence of \(|O|^2\) terms, either ab initio or induced when \( H \) gets an expectation value, may cause one or more of the \( \phi_i \) (and thus \( O_i \)) to develop an
expectation value $w_i$. This gives a mass $m_Y \sim g_{hi} w_i$ to some or all of the hidden gauge bosons $Y$. Let us assume, purely for simplicity, that they are all massive.

The nonzero $w_i$, aside from their effect through $fH \dagger H \langle O \rangle$ on the Higgs boson mass, also cause mixing of $H$ and $\phi$ itself. Said another way, the operator $O$, after symmetry breaking, is now just an ordinary particle at leading order

$$O_i^j = w_i^j w_j + v^j \delta \phi_i \dagger + w^j \delta \phi_i + \delta \phi_i \dagger \delta \phi_j (1)$$

and thus the mixing between $h$ and $O$ is ordinary particle mixing between $h$ and $\phi$. Therefore the unparticle propagator develops poles at $m$ and at $m_h$, with non-zero imaginary parts, and with residues which depend on $f$, $v$ and $w_i$. If the $\phi_i$ have different masses, or the $w_i$ are not all equal, there will be a separate pole for each $i$.

This has the effect that $O$ (or $\delta \phi$, to be less obscure) has an amplitude to decay by $\delta \phi_i \to h^*$ to any kinematically allowed final state of the Higgs boson, such as $b\bar{b}$, $\tau^+\tau^-$, $WW$, $ZZ$, etc. This is precisely what happens in models \cite{18, 19} where there are standard-model singlets that can mix with the Higgs boson. This is also what can happen in hidden valley models \cite{1, 2}, which easily produce sterile scalars that can mix with the Higgs.

But there is another important effect to consider. Just as the Higgs boson can decay to $WW$ and $ZZ$ because it gives them their masses, the $\delta \phi_i$ can decay, if $w_i$ is nonzero and the $Y$ is sufficiently light, to $YY$. And since the Higgs can mix with the $\delta \phi_i$, it too can decay by $h \to YY$.

More precisely, the Higgs and the $\phi_i$ form a mixed system of scalar mass eigenstates $\phi_a$, each of which can decay to standard model fermions, standard model gauge bosons, and hidden gauge bosons, if kinematically allowed. Moreover, they can decay to each other, because of the $H^i H^j \phi_i \phi_j$ coupling, which after symmetry breaking induces many three-particle couplings, such as $h \to \phi_i \phi_j$ and $\phi_i \to hh$, or more generally $\phi_a \to \phi_b \phi_c$. (There are also in general pseudoscalar states as well, but to keep the discussion under control we ignore them here; a more serious study might reveal additional signals.) Whether these three-point couplings dominate over the decays to standard model particles or to $YY$ depends strongly on the couplings and the masses and the mixing angles of all the states.

Moreover, the $Y$, now massive, may not be invisible. It will mix with the $Z$ boson, with a model-dependent mixing angle, and decay with a model-dependent lifetime to standard model fermions: quark pairs or lepton pairs. Its mass is often small and its mixing must be small, for consistency with direct and indirect LEP bounds, so it may well decay with a displaced vertex. For instance, if the $Y$ field is heavier than $2m_h$, then its lifetime is roughly of order that of the $Z$ (about $10^{-24}$ seconds), divided by the square of a mixing angle, times $(m_Z/m_Y)^5$ \cite{1}. For lighter $m_Y$ some decay channels are kinematically forbidden and the lifetime becomes longer. For a 20 GeV $m_Y$ to decay inside the detector requires a mixing angle larger than $10^{-6}$, which is certainly permitted by experimental constraints.

In \cite{14}, where the unparticle developed an expectation value, these issues were not considered. Admittedly it is difficult to express or even recognize these phenomena in purely unparticle language. It is tricky, at best, to write a shift from one vacuum to another, and the Higgs mechanism, in terms of gauge-invariant local operators. For example, unparticle interactions clearly cannot be neglected when computing the two-point function of an unparticle in a shifted vacuum, since three- and higher-point functions in the unshifted vacuum will contribute to the two-point function in the new vacuum. As all known nontrivial conformal field theories in four dimensions are gauge theories, an unparticle’s interactions with other unparticles that contain the hidden gauge fields $Y$, such as $\phi_i D_a \phi_j$ in our toy models, must be considered. And in the toy models above, three-point functions, which contribute to decays, get new contributions from the unparticle interactions and from Higgs-Higgs-unparticle couplings.

1. Possible Signatures

With all of the above effects accounted for, the potential for striking phenomenology emerges. Instead of an invisible sector with a cross-section given by one or another of Figs. \cite{10, 3} or in the figures of \cite{14}, we may have a flurry, or perhaps even a blizzard, of visible final states. Several of these are illustrated in Figs. \cite{10, 11} \cite{12, 13, 14} Decays of any of the scalars $\phi_a$ (the mass eigenstates which are mixtures of the Higgs and the $\phi_i$) may generate final states that range from the well-known to the exceptional. Examples (with resonant pairs shown in brackets) include

- $gg \to \phi_a \to \phi_b \phi_b \to (b\bar{b})(b\bar{b})$ (Fig. \cite{10})
- $gg \to \phi_a \to YY \to (\ell^+ \ell^-)(\ell^+ \ell^-)$ (Fig. \cite{11})
- $gg \to \phi_a \to YY \to (q\bar{q})(\ell^+ \ell^-)$
- $gg \to \phi_a \to \phi_b \phi_b \to (YY)(YY) \to [(q\bar{q})(\ell^+ \ell^-)][(\nu\bar{\nu})(\ell^+ \ell^-)]$ (Fig. \cite{12})
- $gg \to \phi_a \to \phi_b \phi_c \to (YY)(\phi_d \phi_d) \to [(q\bar{q})(\ell^+ \ell^-)][(b\bar{b})]$ (Fig. \cite{13})
- $gg \to \phi_a \to \phi_b \phi_c \to (YY)(\phi_d \phi_d) \to [(q\bar{q})(\ell^+ \ell^-)][(b\bar{b})(b\bar{b})]$ (Fig. \cite{14})

Clearly this is not the entire list. These decays are in addition to classic decay modes, such as $\phi_a \to W^+ W^-$ or $b\bar{b}$, if kinematically allowed. Note the decays are mainly on shell; the fermions in the final states form resonances pairwise, so plots of the invariant mass of the dileptons will show peaks, making the four-lepton channel completely spectacular. (As this paper was completed, I learned that model A, investigated already in
is being further studied in [57], where some of these decay modes of the Higgs boson are also noted.) Displaced vertices from $Y$ decays (or perhaps even from $\hat{\phi}$ decays) may also be present, adding additional spice to the story and reducing backgrounds.

**FIG. 10:** The Higgs (or any sufficiently heavy $\hat{\phi}^o$) may decay to two $\phi$ particles, which each decay to heavy flavor; see [19] for similar examples.

**FIG. 11:** The Higgs (or any sufficiently heavy $\hat{\phi}^o$) may decay to two $Y$ bosons, which decay to two standard model fermions each, often resulting in two-lepton-two-jet or four-lepton final states.

**FIG. 12:** The Higgs (or any sufficiently heavy $\hat{\phi}^o$) may decay to two $\phi$ particles, which each decay to heavy flavor; see [19] for similar examples.

Even this is not all. If the $Y$ is rather light, but the hidden gauge coupling is not small, then it is easy for one or more $Y$ particles to be radiated in the production process, as in Fig. 3. These can then decay to additional standard model particle pairs, as in Fig. 14 again perhaps with a displaced vertex. And still more complexity may arise in models B and C if the nonabelian $Y$ fields, whose masses are related to the $w_1$, can decay to one another, increasing the multiplicity of final-state particles still further.

We should not forget the unparticle production cross-section near threshold for the lightest $\phi$. We still have a continuum above $2m$ and $\phi$-onium resonances just below $2m$. But the number of resonances depends on $m_Y$; as the Compton wavelength of $Y$ decreases, so does the number of resonances, eventually to zero. As the resonances decay, or annihilate, they may produce mainly $Y$ bosons, whose decays would make for striking events. Or if the $Y$ is too heavy, decay may occur through an off-shell Higgs boson. However, the rate for this resonance production may be very small.

By now the fact that this is a hidden valley model should be fully clear. We have new light neutral resonances, with long lifetimes, appearing in potentially large numbers through decay cascades and through radiation [1]. There are new Higgs decays [1, 2]. If this is a supersymmetric world, as in model C, we will also have the physics of [3], whose details depend on the relative masses of the standard model and hidden sector LSPs. Amid all of this, the measurement of the conformally invariant high-energy tail on the production process may seem less urgent, though it remains a very important probe of the hidden sector. Whether it is easy or difficult to measure this tail clearly depends on the visible signal, which may either brightly illuminate or badly obscure the kinematic}

many resonances, has also been considered [18]. But the fact that multiple cascade decays can so generically lead to multiple resonances and high multiplicity in the final state has only been recently emphasized [1, 18] and the possible relevance of displaced vertices for discovering the Higgs boson has apparently also been overlooked until recently [1, 2, 20] (see also comments in [16].)
heavy mesons \( \sim \chi \) rename the light ones.

ple version of this case, with fermions instead of scalars whose mass operator has the lowest dimension.) A sim-

ples most strongly at low energy only to those scalars sector break the flavor symmetry, so that the Higgs cou-

natural scenario is that other couplings in the hidden

k

plings to the hidden sector break its

standard model fermion, \( u, d, s, c, b, t, e, \mu, \tau, \nu \).

dimension.

variable whose behavior is controlled by the “unparticle” dimension.

F. The Unbroken Phase: A Second Look

Now we should revisit the unbroken phase, to see what happens to the massless hidden gluons in models B and C. Of course the unbroken phase is a confining phase at low energy. If the conformal heavy gauge coupling is weak, then the confinement scale \( \Lambda \) is far below \( M \), and any hidden-sector hadrons may be invisible. But if the gauge coupling at the fixed point is fairly strong, as is usually the case when there are large anomalous dimensions for any operators, then confinement may kick in within an order of magnitude or two of the scale \( M \); see for example [29, 37]. In this case, we will see the remarkable phenomenology of a confining hidden valley model, a few examples of which were given in [1].

There are a number of different possibilities, depending on the precise nature of the matter in the hidden sector and the masses of the matter. We will cover just a couple of them; this is not the full range of possibilities!

1. Light Matter in the Fundamental Representation

Suppose first that \( N_f > 1 \), and that the Higgs couplings to the hidden sector break its \( U(N_f) \) symmetry, so that \( k \) of the \( \phi \) fields end up heavy, while \( N_f - k \) are light. (As stated, this is a bit fine-tuned. A more natural scenario is that other couplings in the hidden sector break the flavor symmetry, so that the Higgs cou-

ples most strongly at low energy only to those scalars whose mass operator has the lowest dimension.) A sim-

ple version of this case, with fermions instead of scalars and \( N_f = 2, k = 1 \), was discussed in [1].

We will continue to call the heavy fields \( \phi_i \), but will rename the light ones \( \chi_r, r = 1, \ldots, N_f - k \). The heavy-

heavy mesons \( \sim \phi^i \phi^j \) will be rarely produced, except for \( \phi \)-onium states, which we will ignore until the next section. But the heavy-light \( \phi \chi \) mesons, analogous to \( B \) and \( D \) mesons, will always be produced in open \( \phi \) production, through the process \( gg \rightarrow h \rightarrow \phi \phi \). (Here the Higgs boson may be on- or off-shell.) The light-light \( \chi \chi \) mesons, analogous to \( \rho, K \) and \( \pi \) mesons, will also be produced in open \( \phi \) production. Thus, just like open-charm production in the standard model, \( \phi \phi \) pair production typically leads to two stable (and invisible) heavy-light mesons, along with some number of light mesons. The higher the center-of-mass energy, the more light mesons produced. At high enough energy this process is driven by a parton-shower of hidden sector \( Y \) bosons.

The visibility of the signal then depends on whether there are \( \phi_i \rightarrow h \phi_j \) decays, as discussed in Sec. [11] and on whether the light mesons can decay to the standard model. By assumption the Higgs has smaller couplings to the \( \chi_r \) than to the \( \phi_i \), but they need not be zero, so the Higgs can mediate decays of some light meson states. Other small couplings, possibly higher dimension operators that couple hidden-sector operators involving \( \chi \) to standard model operators, may also mediate light meson decays. One example was given in [1]; many other examples, with varying phenomenology, may be invented. If one allows oneself the freedom to introduce all possible operators, as in [31], then essentially any decay modes for the light mesons allowed by symmetry, with a vast range of lifetimes, are possible.

Thus, the signatures in models B and C may include the following (see Fig. [15]):

- For each of the signatures outlined in Sec. [11] and the invisible production processes in that section, we should add one or more light mesons.

- As illustrated in [1], these mesons are most likely to be scalars and pseudoscalars decaying to heavy quarks and leptons, or to gluon/photon pairs, and possibly spin-one mesons decaying to leptons and quarks more democratically.

- Displaced decays are possible for light pseudo-

scalars and very light vectors, or for any light state if the \( \chi_r \) couple much more weakly than \( \phi_i \) to the standard model.

Other hidden hadrons with more complicated decays are certainly possible [1]. For instance, in model C, we also have supersymmetric partners of these states, which if supersymmetry is not badly broken may allow for fermionic hadrons with additional three-body decays. Decays to four particles are also not uncommon. Indeed there is much more to say about the supersymmetric case, and the interplay of supersymmetry breaking with hidden-valley phenomenology.

Note that from the unparticle point of view, the pro-

duction of \( n \) light mesons and two heavy mesons must be described using an \( n + 3 \)-point function involving gauge invariant operators built from \( \chi \) and \( \phi_i \). Again, this correlation function would be highly suppressed in the
In a confining theory with light propagator: in such a theory, the propagator is not a cutily break. This has a dramatic effect on the unparticle one another; they are bound by a string that cannot eas-
in only at continuing up to very high energy, with a cut setting
space for decays of this class. Without including unpar-
with relatively narrow widths, along with the large phase
arise from the production of many light on-shell states
conformal limit compared to other processes, but here it
benefits from the cuts and resonant enhancements that
arise from the production of many light on-shell states
with relatively narrow widths, along with the large phase
space for decays of this class. Without including unpar-
ticle interactions and treating them with great care, one
might overlook these complex yet dominant processes.

2. Heavy Matter Only

Now suppose instead that all the matter in the hid-
den sector becomes massive relative to Λ after conformal
symmetry breaking. Then the low-energy limit is a pure
hidden Yang-Mills theory, with a hidden confining flux
tube. This very generic situation was also considered in
\[1\]. The low-lying hidden hadrons are hidden glueball
states, of mass \(\sim \Lambda\).

When a pair of \(\phi_i\) is produced, they cannot escape
one another; they are bound by a string that cannot eas-
ibly break. This has a dramatic effect on the unparticle
propagator: in such a theory, the propagator is not a cut
above \(2m\) but is instead a sum of bound state resonances
continuing up to very high energy, with a cut setting
in only at \(q^2 = (4m)^2\), where two such resonances can
be pair produced. Note these states are not stable: all
can decay by emission of hidden glueballs except for the
lightest, which can decay via annihilation to two or more
hidden glueballs. The widths of the resonances grow as
one goes up the tower, just due to the phase space for the
decays, so eventually they blur together to form a con-
tinuum which must satisfy the contraints of conformal
invariance. Where this happens depends on the details
of the theory, but above this scale the calculations done
with conformal unparticle propagators will apply.

As an aside, note that in the case where the fields
\(\phi\) carry electric charge or color, they are often called
"quirks" \[1, 39\]. In this case their bound states may
also radiate standard model gauge fields. This was very
briefly considered in \[1\] for confinement scales well above
1 GeV, but the physics is very subtle, and the question of
the expected LHC signals is still under study \[40\]. The
case of quirks with confinement scales well below 1 GeV
has many special features and has been considered by
\[39, 40\].

Clearly the observability of the physics depends in this
case on whether the hidden glueballs can decay visibly.
According to lattice simulations \[21\] there will be numer-
ous glueballs, of various spins, parity and charge con-
jugation, which cannot decay to other glueballs. How
each one decays — its decay mode, lifetime, and branch-
ing fractions — depends in detail on whether there are
nonzero couplings between hidden gauge invariant opera-
tors \(\text{tr} F_{\mu\nu} F^{\mu\nu}, \text{tr} F_{\mu\nu} F^{\mu\nu}, \text{tr} F_{\mu\nu} F^{\rho\nu}\), etc., and operators
in the standard model. Because of the high dimension
of these operators it is easy to arrange that all these de-
cays would occur outside LHC detectors. But if one or
more of the lifetimes is sufficiently short, the result will
be a classic signature of a hidden valley: long-lived light
neutral resonances produced in abundance, with likely
missing energy and large event-to-event fluctuations.

In particular, hidden \(\phi\)-onium production will produce
some number of low-\(p_T\) glueballs, with various \(J^{PC}\) quan-
tum numbers, emitted as a \(\phi\)-onium state relaxes to the
ground state, followed by a blast of glueballs produced
in the annihilation process; see Fig. 16. The signatures
are then

- A number of high-\(p_T\) glueballs with \(p_T \sim m\), and a
number of low-\(p_T\) glueballs with \(p_T \ll m\).
- Decays of the various glueballs, whose masses span
a range of about a factor of three, to different final
states, such as \(gg, b\bar{b}, \gamma\gamma, ggg\), etc.
- Possible displaced vertices from one or more glue-
ball decays.
- The \(\phi\)-onium annihilation may occasionally occur
through \(h^*\); producing any final state of the Higgs,
such as \(ZZ\), in place of the high-\(p_T\) glueballs.

Unfortunately there is no known method for obtaining
reliable predictions for the \(p_T\) spectrum of the glueballs
and their standard model daughters. Much additional
work on this example is needed.

3. Possibilities for Future Study

There are many other possibilities. The hidden gauge
group may not be SU\((N)\); either the \(\phi\) fields or the light
\(\chi\) fields, if any, may not be in the fundamental represen-
tation; the phase of the theory may be partially broken
and partially confined. Each of these possibilities will
change the details drastically, but in the majority of cases
the basic features of the hidden valley scenario, and its
experimental implications, will be retained. While the

\(\chi\) pairs

\(F\) \(\mu\nu\), \(T\) \(\mu\nu\), \(p\) \(\mu\nu\), etc., and operators
\(\text{tr} F_{\mu\nu} F^{\mu\nu}, \text{tr} F_{\mu\nu} F^{\mu\nu}, \text{tr} F_{\mu\nu} F^{\rho\nu}\), etc., and operators
\(\text{tr} F_{\mu\nu} F^{\rho\nu}\), etc.

\(\phi\)-onium annihilation may occasionally occur
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\(\chi\) pairs
conformal invariance of the theory and the dimensions of the operators do constrain some inclusive observables, the most dramatic effects on the phenomenology are beyond the easy reach of unparticle methods.

G. Other Models, Other Couplings

We have seen that a wide variety of signals can arise even in simple toy models. There is an enormous diversity of phenomenological possibilities, as is typical in the hidden-valley scenario. But we have only discussed a very small set of unparticle models, those with a scalar unparticle with dimension somewhat below 2, with a coupling to the Higgs boson. Are these cases special? Also, we have often used weak-coupling intuition for guidance. Is this misleading?

1. Stronger coupling and $d_\mathcal{O} \ll 2$

In model C, as we decrease $N_f/N_e$, it is known that the hidden gauge coupling becomes stronger and $d_\mathcal{O}$ decreases from 2 toward 1. Should we expect something completely different from our discussion above as $d_\mathcal{O}$ approaches 1?

The reader is invited to consider the case $d_\mathcal{O} = 1 + \epsilon$. There the operator $\mathcal{O}$ can be treated a scalar that couples weakly to a conformal sector, for example through a linear coupling to an operator of dimension $3 - \epsilon$, or quadratically to an operator of dimension slightly less than 2. The arguments can be repeated using toy models. Excellent toy models are Banks-Zaks supersymmetric fixed points with additional gauge-singlet superfields, as in Seiberg “magnetic” fixed points, coupling in the superpotential to squark-antisquark bilinears. Not surprisingly, since this is just model C with additional scalars, and since we already considered model C with additional scalars (the Higgs boson itself!) in our examples above, the physics in this case has the same features, somewhat rearranged. The resonances (both strong and weak), mixing, cascade decays and complicated multibody final states that we have seen here may arise there as well, though with rates that decrease to zero as $\epsilon \to 0$. The exercise is left for the reader.

2. If there is no coupling to the Higgs boson

If the $|H|^2 \mathcal{O}$ coupling is absent, then some of the effects discussed in the earlier sections may be absent as well. In particular, we may not see unusual Higgs boson decays, or the effect of Higgs mixing with the unparticle sector. However, we may still see unusual decays of other heavy particles, including very rare decays of $Z$, $W$ or $t$, or very common decays of new particles, such as supersymmetric partners, little Higgs partners, Kaluza-Klein partners, $Z'$ bosons, right-handed neutrinos, or new scalars other than the Higgs boson. What matters more than the Higgs coupling is whether there is a mass gap which is sufficiently large.

How might a large mass gap (larger than a few hundred MeV, at least) naturally arise without the Higgs boson coupling? This is very easy to imagine, since we know there must exist some mechanism that generates the electroweak scale in the standard model sector. For instance, supersymmetry breaking in our own sector, if generated via gauge mediation or supergravity mediation, will naturally generate supersymmetry breaking in other sectors as well. While this breaking may be somewhat suppressed, it may still lead to a mass gap (or ledge) at or somewhat below the 100 GeV scale. Technicolor models may easily break symmetry groups larger than that of the standard model, including those of hidden sectors. A further possibility is that the very mechanism that leads to the local couplings at low energy between the two sectors is precisely the same as that which generates the mass gap: for example, if there are massive particles charged under both groups, the masses at 1 to 10 TeV may destabilize a quasi-fixed-point, causing the hidden gauge coupling to run strong and confine at a scale at 1 to 100 GeV. The point is that it is very easy to imagine models with a large mass gap without invoking the coupling to the Higgs boson.

The physics that ensues can, as before, leave the theory in any number of phases: broken, confined, or partly both, with any number of possible light modes that can decay via couplings to the standard model. The question of whether the hidden sector is invisible is very model-dependent.
3. Effects in top quark decays

If there are couplings to the top quark, then even without couplings to the Higgs boson one may find remarkable signals. For example, in the $t \to cO$ transition considered in [5], one may have, instead of missing energy, a decay to many particles, such as in Fig. 17 where the top quark decays to nine particles, including four hidden-sector resonances decaying to two particles each. This particular process can occur if the hidden sector is in a broken phase, where the $\phi \to YY$ decay can occur. Note this is a fully reconstructable top quark decay.

The total differential rate for this process may still be given by the unparticle prediction [5], because the energy released in the decay $t \to cO$ may be large, compared to the masses of the $\phi$ particles. But unfortunately the kinematic variable which one needs in order to measure the unparticle dimension involves identifying which of the jets is the charm quark. This may in some cases be the most energetic jet, but clearly it will be very challenging to make the inclusive measurement and determine the unparticle dimension directly. A similar problem would arise in other visible signals, such as $t \to cbbb$, or $t \to c\gamma\gamma$, with or without missing energy. However, this problem may be absent if the hidden sector particles (such as the $Y$ in Fig. 17) decay with a displaced vertex.

For instance, a neutralino – perhaps, in our toy model C, through its Higgsino coupling to the unsparticle – may decay to $\phi$ and its superpartner, which in turn decays to $\phi$ and $\tilde{Y}$, the invisible hidden gaugino (Fig. 18). This type of decay can significantly reduce the missing energy signal which is typically used to find supersymmetry, and can replace it with soft jets and leptons [3]. More research to understand how to find such a signal is needed.

Another interesting possibility [3] is that the lightest standard model superpartner is not neutral. For example, if this lightest standard model particle is a stau, then the decay $\tilde{\tau} \to \tau \tilde{O}$ to an unsparticle is similar to the $t$ decay of [3], in that the $\tilde{\tau} \to \tau$ kinematics may reflect the dimension of $\tilde{O}$. However, as with the top quark decay mentioned above, the visibility of the kinematical power law will depend on the details of the final state emerging from the unsparticle, as it is converted into visible particles. Note also that the $\tilde{\tau}$ may easily be long-lived (this is less likely in the neutralino case) and may decay with a displaced vertex [3].

Essentially the same physical phenomena can arise in any model with new particles that carry a new conserved global symmetry, such as KK-parity in extra-dimensional models, or T-parity in little-Higgs models. As long as there are particles in our sector and in the hidden sector carrying the new charge, the possibility of interesting cross-sector decays exists. The scenario of a new heavy particle in the standard model (the lightest particle carrying the new charge) and an unsparticle sector with conformal invariance and a mass gap below 100 GeV (which will have a lighter particle carrying the new charge) virtually guarantees the decays studied in [3] will occur. The only question is whether the decays are visible, and this depends on the details of the hidden sector and the size of the mass gap.

4. Supersymmetric decays, and analogous cases

With supersymmetry, unparticle couplings to supersymmetric particles can lead to challenging decays of the lightest standard model superpartner, as discussed in [3].

Obviously there is much more to do in this arena. Ex-
III. THE STEPHANOV MODEL AND HIDDEN VALLEYS

An interesting approach to unparticles was provided by Stephanov [15], where an unparticle was modeled by appealing to the five-dimensional language inspired by the AdS/CFT correspondence, also known as “gauge/string”, “gauge/gravity”, or “boundary/bulk” duality [22, 23, 24]. (See also the work of Randall and Sundrum [8, 9].) The five-dimensional language, as we will see, is indeed instructive, but only will guide us if we use it fully. We will need the full gauge-string correspondence, not the watered-down gauge/gravity version, in order to capture all the physics and see the hidden valley in full (s)unshine.

A. The Stephanov Viewpoint and Hidden Valleys

In order to regulate and then interpret the unparticle propagator, Stephanov broke the conformal invariance with an infrared cutoff. (Note this is not deconstruction, but compactification [62, 63].) This cutoff is often used in the gauge/string literature as a model of how to describe confinement in the gauge theory in terms of a five-dimensional (5d) theory on a warped space-time. In short, Stephanov’s method of regulating the unparticle propagator is simply this: a model of a two-point function in a hidden sector which is confining in the infrared and conformal at larger energies. Many examples of such theories are known; one explicit example with a dual string description is given in [29].

Models of the same type were already considered in Sec. II F. But the predictions presented in that section are very different from those of Stephanov. This makes it far from obvious that we are dealing with a hidden valley. Let us review these predictions and see where the difficulties with them lie.

First, Stephanov predicted a narrow tower of states. Second, he suggested these states would decay to standard model particles through the coupling to the unparticle. Third, he suggested these states could have very long lifetimes and could be detected from displaced vertices. Finally, he suggested the lifetime \( \tau_n \) of the \( n^{th} \) state would be related to the dimension of the unparticle by \( \tau_n \sim m_n^{-2d_s} \), where \( p \) is a (known) positive number.

This tower of long-lived states, with a directly measurable lifetime-to-mass relation set by the dimension of the unparticle, is an impressive prediction of unparticles not shared by typical hidden valley models. But there is a good reason for this.

The point is that Stephanov’s formulas apply for a hidden gauge theory in the limit that the number of colors \( N \), and the ’t Hooft coupling \( \lambda \equiv \alpha N \), where \( \alpha = g^2/4\pi \) is the hidden gauge coupling, are extremely large — how large will be explored below. (Meanwhile \( N_f \), the number of flavors of matter fields, must remain finite and small, so these are not classic Banks-Zaks-type or Seiberg-type fixed point theories.) In this limit, a confining gauge theory has a spectrum that is merely an infinite tower of stable non-interacting hadrons of spin \( \leq 2 \). Many familiar aspects of gauge theory would be absent, including all high-spin hadrons, BFKL dynamics, parton showers, and the like. This differs so dramatically from QCD, and even from gauge theories with dozens of colors, that it was not recognized as a hidden valley. But the interpretation is this: if one were to take a hidden valley model into this extreme regime, it would eventually resemble the narrow-tower model. Less extraordinary models will have, not surprisingly, more ordinary predictions similar to those of [1].

B. Stephanov’s approach

In a companion paper [61], I will more carefully and pedagogically add in the \( 1/N \) and \( 1/\sqrt{\lambda} \) corrections to the “narrow-tower” model. I will also consider some variants of this model and see how easily the predictions may be altered. (For instance, in many realistic models, the tower may have a finite or infinite number of states extending only over a finite range of energies, with a continuum above. The continuum may also have additional embedded resonances. Alternatively, the density of narrow states may suddenly change by a finite factor at some scale. The spacing of the states need not be uniform; there may be degeneracies of states that grow with \( n \); etc.) I will also consider further how hard one must push the theory to make the original narrow-tower model of [15] appropriate to the physics. Here, I will keep things short and state some of the more subtle claims without proof.

First some notation. The radial coordinate in \( AdS_5 \) will be called \( r \), running from a boundary at \( r = \infty \) to a horizon at \( r = 0 \); the five-dimensional metric is

\[
ds^2 = \frac{r^2}{R^2} (-dt^2 + dx^2 + dy^2 + dz^2) + \frac{dr^2}{R^2}.
\]

For RS experts, my coordinate is chosen such that if I cut off the space at \( r = r_{UV} \) for large \( r \) and at \( r_{IR} \) at small \( r \), then in RS1 the Planck brane is at \( r_{UV} \) and the TeV brane at \( r_{IR} \). In RS2 the interpretation would be slightly different; but in any case, the space represents in this case a theory which is conformal between the two energy scales \( \mu_{UV} \sim r_{UV}/R^2 \) and \( \mu_{IR} \sim r_{IR}/R^2 \). With no UV or IR cutoff, and with the addition of a well-behaved five dimensional compact space \( X \) to make a total of ten dimensions, a superstring theory on this space would, according to gauge/string duality, precisely represent a conformal gauge theory.
Stephanov considered an AdS space with $r_{UV} = \infty$ and $r_{IR}$ finite. In a fully consistent setting, this would correspond to a theory which is conformal at high energy but at low energy has some sort of conformal symmetry breaking. Cutting off the space sharply, without any nuances, is called the “hard-wall” model. It has been used extensively for study of confining gauge theories at large 't Hooft coupling in gauge-string duality \cite{41,42,49,50,51,52}. It has its limitations, but is often useful. In such a model, the scale $\mu_{IR} \sim r_{IR}/\Lambda_{c}$ is of order $\Lambda$, the confinement scale. Other models \cite{43} give similar structure, though the details differ.

It is a natural conjecture (independent of $\lambda$) that in the $N \to \infty$ limit of a conforming gauge theory the two-point function of a reasonable operator can be exactly written as

$$\langle O(q)O(-q) \rangle = \sum_{n} \frac{|F_{n}|^{2}}{q^{2} - m_{n}^{2} + i\epsilon}$$. (3)

where the masses $m_{n}$ are those of the confined hadrons $|n\rangle$ created by acting with the operator $O$ on the vacuum,

$$O|0\rangle = \sum_{n} F_{n}|n\rangle$$. (4)

Within the hard-wall model and its cousins, and using the low-energy five-dimensional gravity theory, these equations can be shown to be true without subtleties. Thus this equation (and its analogues for higher spin) is correct for $\lambda \to \infty$, $N \to \infty$, $\Lambda/N$ fixed and not too large, at least for primary operators with spin and dimension of order 1.

One can see that within the hard-wall model, and many of its variants, the only significant change as $d_{O}$ changes is the $n$-scaling of the $F_{n}$, as reviewed in \cite{15}. The $m_{n}$ change; their $n$-scaling typically does not. However, this precise feature is a property of a particular model, and as noted in \cite{15} the constraints are rather weak.

Conformal invariance only requires that the two-point function approaches a particular power law. In the limit that $r_{IR} \to 0$, as reviewed in \cite{15}, the spacing between the modes goes to zero, and the two-point function must regain its conformal form. But this requirement imposes only a single relation between the large-$n$ behavior of the $m_{n}$ and that of the $F_{n}$. As emphasized in \cite{15}, the requirements of conformal invariance on Eq. (3) do not permit the mass spectrum of the tower to be predicted from the operator dimensions alone. Thus to measure $d_{O}$, one must measure something else.

C. Finite $N$ Effects on the Spectrum

It \cite{15} it was proposed that one should measure the lifetimes of the states. Let us review the calculation. Although there are technical problems with the particular case chosen, the basic logic is correct: with a coupling in the Lagrangian $\sim \hat{c} O_{SM} O$, where $O_{SM}$ is a standard model operator, the decay of a state $|n\rangle$ in the tower to a standard model state such as $\mu^{+}\mu^{-}$ is proportional to

$$|\hat{c}\langle \mu^{+}\mu^{-}|O_{SM}|0\rangle(0)|O|n\rangle|^2$$ (5)

For example, in the case \cite{15} considered, one obtains (converting to a notation in which $m_{n} \sim \Lambda\sigma$ at large $n$) a decay rate of the form,

$$\Gamma_{|n\rangle-\mu^{+}\mu^{-}} \sim \alpha_{\text{eff}}^{(n)} m_{n}$$ (6)

where the effective coupling is

$$\alpha_{\text{eff}}^{(n)} = \frac{c^{2}A_{O}}{16\pi^{2}} \left(\frac{m_{n}}{M_{Z}}\right)^{2}(d_{O}-1)$$ (7)

Here $A_{O}$ is a normalization constant for the unparticle, of order 1, and we have converted $\hat{c}$ into a dimensionless constant $c$ times the appropriate power of $M_{Z}$, following \cite{15}. This rate can be enormously suppressed if either (1) $c$ is very small, or (2) $\Lambda, m_{n} \ll M_{Z}$ and $d_{O}$ is significantly above 1.

We see, therefore, that if the lifetimes of the states can be measured, then so can $d_{O}$. But how can they be measured?

1. A necessary condition

Can these states decay with visibly displaced vertices? This requires lifetimes in the picosecond range or longer. For decays to electrons (muons), the electron (muon) mass is of order 1 (100) MeV, so let us take $m_{n} \sim 1$ (100) MeV as well, to lengthen the lifetime as much as possible. Then we must have

$$\alpha_{\text{eff}}^{(n)} \lesssim 10^{-9} \quad \text{(decay to electrons)}$$ (8)

$$\alpha_{\text{eff}}^{(n)} \lesssim 10^{-11} \quad \text{(decay to muons)}$$ (9)

These tiny numbers are already an issue since production rates, compared to ordinary electromagnetic processes, are very small. However, there is a more serious issue.

2. Why the previous condition is not sufficient

The above condition for displaced vertices, while necessary, is not sufficient! It is only appropriate if the state $|n\rangle$ has no other decay modes. And this is not true, except typically for the ground state and perhaps the first excited state in every tower.

The narrow-tower model implicitly assumes that the 5d scalar field that represents the scalar unparticle is non-interacting. This is equivalent to assuming that the unparticle has no 3- or higher-point functions. This is true in the $N \to \infty$ limit (and also in the free $d_{O} \to d_{\text{min}}$ limit, where $O$ becomes an ordinary free
particle. But the infinite $N$ limit is very misleading (as is the $d_\text{O} \to d_{\text{min}}$ free particle limit.) At any finite $N$ a conformal gauge theory will have $k$-point functions for $k > 2$. Equivalently, the 5d scalar field \textit{will have self-interactions} (if $d_\text{O} > d_{\text{min}}$). As a result, the unparticle now has higher-point functions.

Also, one cannot treat one unparticle tower in isolation. There are always other fields in the bulk. At the very least, the 5d graviton must be present, because it represents the energy-momentum tensor of the hidden sector, which is part of any conformal theory. Each such field, in a confining gauge theory, will have its own tower. And of course the graviton interacts with itself and with all other fields in the bulk. In fact, in any conformal gauge theory one expects many fields in the bulk, with many quantum numbers. Any conserved currents in the theory, for instance, will be represented as 5d massless gauge fields, and they too will interact with themselves, with the graviton, and with any 5d fields that carry the corresponding conserved charge. All this is to say that there is no conformal gauge theory without three-point functions and operator product expansion (OPE) coefficients — except at $N \to \infty$ — and that $T_{\mu \nu}$ and conserved currents $J_\mu$ always have a nontrivial OPE.

Once these interactions are introduced, we no longer expect a tower of extremely narrow states decaying to standard model particles. Any state with high mass will decay via these interactions. The states may still be relatively narrow, but nowhere near as narrow as predicted in \cite{49}.

3. \textit{Another necessary condition}

Let us now estimate the widths of the excited states. (We assume $d_\text{O}$ is not very close to $d_{\text{min}}$; otherwise special treatment is required.) The width of the $n^\text{th}$ state ($n > 1$) to other hidden sector states, may be very roughly estimated as

$$\sum_{n', n''} \Gamma_{n \to n', n''} \sim \frac{g(n) m_n}{8 \pi N^2}$$

where $g(n)$ characterizes the growth in the number of decay channels and monotonically grows with $n$. The states become narrow rather quickly with $N$, for fixed $n$, but conversely their widths grow with $n$. Suppose as $n$ becomes large that $m_n \sim n^\sigma \Lambda$ and $g(n) \sim C n^\beta$. Typically $\sigma \leq 1$ (it is $\frac{1}{2}$ in QCD and in string theory and is 1 in many gauge-gravity duality examples.) Meanwhile one might naively expect $\beta \sim 2$, accounting for the scaling of available channels with $n$, but in most computable theories all but $n$ couplings are small, as in \cite{49}, so to be conservative let us only assume $\beta \geq 1$. Finally $C > 1$ accounts for the presence of multiple towers of states in the theory, which provide multiple classes of decay channels. If $N_f$ is large, then $C \sim N_f$; the situation for small $N_f$ is less clear, but $C$ is certainly larger than 1. The states bleed together when

$$m_n - m_{n-1} \sim n^{\sigma - 1} \Lambda$$

is of order

$$\Gamma_n \sim \frac{C n^{\sigma + \beta} \Lambda}{8 \pi N^2}$$

and thus occurs at

$$n \sim \left[ \frac{8 \pi C}{\Lambda N^2} \right]^{1/(\beta + 1)}$$

Since $\beta + 1 \geq 2$, we expect at most the first $N$ states to be narrow relative to their separation. Actually this is often a large overestimate, due to the fact that we have ignored stringy effects; we will return to this in the Sec. III E.

The inverse of these widths puts an upper limit on the lifetimes of the excited states. Unless

$$\frac{g(n)}{8 \pi N^2} \lesssim 10^{-9} \quad \text{decay to electrons} \quad (14)$$

$$\frac{g(n)}{8 \pi N^2} \lesssim 10^{-11} \quad \text{decay to muons} \quad (15)$$

the lifetimes will be too short for displaced vertices. Thus, except for the lowest one or two states, for which $g(n) = 0$, this condition requires $N \sim 10^4$ for electrons and $10^5$ for muons. And this is generous, because we assumed the lowest possible mass for the decaying state. Also, we would hope to see at least four or five states, in order to measure a power law, and $g(n) \sim C n^\beta$ is often large compared to one and grows with $n$.

\begin{itemize}
  \item In short, \textit{we do not expect a tower of states with displaced vertices unless both $N \gg 10^4$ and $\alpha_{\text{eff}}^{(n)}$ is very small.} What, then, is the phenomenology more likely to be?
  \item There are two possibilities, depending on whether the decays of the hidden sector states are to other hidden sector states or to standard model particles. Decays to standard model states will dominate only if
  \begin{align}
    \alpha_{\text{eff}}^{(n)} &> \frac{g(n)}{8 \pi N^2} \\
  \end{align}

  But $\alpha_{\text{eff}}^{(n)}$ cannot be large, or effects from the hidden sector would already have been seen. (In fact, if $\alpha_{\text{eff}}^{(n)}$ is of the same order as $\frac{g(n)}{8 \pi N^2}$, then constraints on $\alpha_{\text{eff}}^{(n)}$ are even stronger than is often realized, because of effects that we will discuss in Sec. III D.) Again, we are forced to take large $N$ — not as large as required for a tower of displaced vertices, but still very large. Being very generous, we would require $N \sim 300$ in almost any conceivable situation; much larger $N$ is required, for instance, in the case of a vector unparticle mixing with the $Z$ boson.
\end{itemize}
4. If decays to standard model particles dominate

If $N$ is of order 100 or more and the unparticle coupling is as large as is allowed by experiment, then there is a narrow window in which

- the widths of the states are determined by the unparticle coupling and
- the widths of the states are large enough to measure.

Assuming resolutions in the few MeV range, and masses in the few GeV range, one might imagine that if $\alpha_{\text{eff}} \gtrsim 10^{-3}$ then one could measure the scale-invariant prediction of [15] through a tower of states with growing widths.

However, if $\alpha_{\text{eff}}$ is too small, the states simply won’t have a measurable width, making the prediction untestable. In this case one can only measure the masses $m_n$. But recall that these are not determined by scale invariance, and cannot be used to measure $d\sigma$. One might hope that since of order $N$ states may have narrow widths, higher states in the tower might always have $\alpha^{(n)}_{\text{eff}}$ so large that the widths are large enough to measure. But a little thought shows there is no guarantee of such a regime. The higher states move closer together as $n$ increases, if $\sigma < 1$, so instead of the widths growing to measurable size, the distance between adjacent states may shrink to unmeasurable size. Also, as $n$ increases so does $g(n)$, so it is possible that the higher states do not decay preferentially to the standard model. And worse, the logic used in the estimates in this entire section breaks down when $n$ is large enough that stringy effects must be accounted for; see Sec. III E.

5. If decays to standard model particle do not dominate

If $N$ is less than 100 or so, or if $\alpha^{(n)}_{\text{eff}}$ is small, the decays within the hidden sector dominate. In this case the lifetimes are not determined by scale invariance; they are determined by $g(n)$, which depends on the details of the hidden sector. Moreover, the partial widths to lepton pairs are typically very small, unless $N$ is very large, making the line-shape of the resonances difficult to observe. This is by far the most likely scenario for a hidden sector!

Thus we are not very likely to observe dilepton pairs from a tower of states. For reasonable values of $N$, and for $\Lambda > 2m_e$, we will find at most $N$ rather narrow states, decaying too rapidly to other hidden sector states for a displaced vertex, and with tiny branching fractions to dileptons. The lifetimes of the states will not be set by the dimension of the unparticle. The lightest state or states in some towers, which are the only states that cannot decay within the hidden sector, will have much longer lifetimes, and may decay to standard model particles with displaced vertices (though possibly outside our detectors.) Only in a narrowly tuned case — small $N_f$, very large $N$, and a coupling large enough that the widths of the first few states are rather large and are determined by the unparticle coupling, might the lifetimes be both measurable and determined by scale invariance.

6. Summary

Let us summarize the implications of this section. Reference [15] includes a correct computation of the partial width for each state in the tower to decay to standard model particles and become observable. But one cannot then assume that this partial width is the total width, and invert it to infer long lifetimes for all the states in the tower. Instead, at reasonable $N$, almost every state in the tower, except the lowest one or two, will decay predominantly and rather rapidly to other states in the tower, or to states in other towers associated to other operators. The lifetimes will be much shorter than estimated in [15], so there will be no displaced vertices. Moreover, the branching fraction to standard model states will be tiny, and it will be very difficult to measure the partial widths. This is simply the statement, familiar from QCD itself, that most hadrons in a gauge theory decay rapidly to other hadrons, and have large widths to do so; their branching fractions to, say, $e^+e^-$ are very small. Thus, unfortunately, the predictions of [15], while possible at extraordinarily large $N$, are unlikely to be seen in nature.

![Cross-Section](image)

FIG. 19: The total cross section $\sigma(s)$ for hidden-sector production in an $e^+e^-$ collider, in the narrow tower model of an unparticle.

D. Unparticle Production at Finite $N$

The cross-section in the narrow tower model takes the form of Fig. 19. In Fig. 20 are shown possible cross-sections for more realistic towers, for small $N$ and large $N$; the small $N$ case resembles low-energy QCD in the $\rho$ channel, and the large $N$ case resembles charmonium production without open charm. But this is the total cross-section for hidden sector production, whereas the cross-section for $e^+e^- \rightarrow \mu^+\mu^-$, both for small and larger $N$, suffers from the low dilepton branching fraction for
FIG. 20: The total cross section $\sigma(s)$ for hidden-sector production in an $e^+e^-$ collider. The thick curve is for small $N$ and resembles QCD; the thin curve is for larger $N$ and more resonances are visible. The first resonance is much narrower than the others, as it can decay only by emission of standard model particles; the others decay within the hidden sector. See [4] for a study of a hidden valley with a light dilepton resonance.

hidden-sector states; it is shown in Fig. 21. Only the first resonance is potentially observable, and since it is small and very narrow, it may easily have been missed up to now.

FIG. 21: A cartoon of the cross-section for $e^+e^- \rightarrow \mu^+\mu^-$ in a model at small to moderate $N$. The falling standard model production rate is supplemented by a single extremely narrow resonance, the lowest resonance in Fig. 20. Its height has been exaggerated greatly for clarity. None of the other resonances have measurable branching fractions to $\mu^+\mu^-$. Only a handful of visible resonances are expected in a generic hidden valley, at most one or two for each tower.

In fact, QCD is an excellent model for the hidden sector. From the point of view of the leptons of the standard model, with electromagnetism turned off, it is a hidden sector, coupled to leptons only by the Fermi interaction. It is no accident that AdS/QCD methods, placed into the hidden sector and treated with care, reproduce QCD-like physics in their effect on $e^+e^- \rightarrow \mu^+\mu^-$. 

Now suppose that we do choose a theory with $N \sim 20$ and choose to run an $e^+e^-$ collider on one of the excited hidden-hadron resonances (say, $n = 10$). What will we see? Certainly we will not observe the process shown in Fig. 22; the branching fraction is too small. Most of the time the resonance will undergo a cascade decay to several light hidden hadrons, each stable against decay to others, which in turn will decay with long lifetimes to standard model particles. This classic hidden valley signature is shown in Fig. 23.

FIG. 22: An allowed but very rare process, in which the tenth resonance is produced by and decays back to standard model fermion pairs.

FIG. 23: A much more common process; the 10th resonance, produced by standard model fermion pairs, decays through a cascade. In the final state appear several particles, each the lightest ($n = 1$) resonance of a tower. Each of these then decays (possibly late) to standard model particles, here assumed to be fermion pairs.

For large $n/N$, the very language that we used in the previous paragraph breaks down. The sum in Eq. (3) is modified in two ways at finite $N$. First, the poles at $m_n^2$ move off the real axis to become resonances. This already gives the two-point function the form shown in Fig. 20. But we also must supplement the formula (3) with cuts from multi-particle production, which is suppressed by factors of $N$ but becomes increasingly important at large $q^2$. We need not produce hadrons only by producing them resonantly, as in $e^+e^- \rightarrow \rho \rightarrow \pi^+\pi^-; \rho$; we may simply produce them directly, as in $e^+e^- \rightarrow \rho^+\rho^-\pi^0$. As $q^2$ increases, the cuts accumulate, eventually stealing support from the resonances and building up the continuum contribution to the two-point function. The dominant production at large $q^2$ is thus not even of single unstable heavy hadrons, but rather of multiple light or moderately light hadrons, whose phase space grows very rapidly as $q^2$ increases. Any excited hadrons among those produced will decay to the lightest ones, and a large number of light hadrons may then decay with long lifetimes to standard model particles. The production process, absent in the
narrow tower model, is illustrated in Fig. 24 of course it also gives a high-multiplicity final state.

Importantly, even if the dilepton branching fractions of the excited states are large, rare high-multiplicity events must not be ignored, as they provide strong constraints on new hidden sectors. Even a small number of events with four or more leptons or photons would have been easily observed, since standard model backgrounds fall so rapidly with multiplicity. In turn, this implies there are much stronger constraints on $\alpha_{\text{eff}}^{(n)}$, and thus on $c$, the unparticle coupling, than have so far appeared in the literature. And in turn, because of Eq. (16), this makes the likelihood of observing excited states with measurably large dilepton branching fractions even smaller, and the likelihood of high-multiplicity states even larger.

![Diagram](image)

FIG. 24: At higher energy, multiple resonances are produced together; each undergoes a cascade decay as in Fig. 23

Thus, instead of a tower of neutral resonances with long lifetimes, the prediction of the narrow-tower model at a reasonable $N$ is indeed that of a typical hidden valley. As in [1], events that access the hidden sector will result in light neutral resonances with long lifetimes, produced in abundance, with large event-to-event fluctuations and possibly large missing energy. Here the abundance arises through the cascade decay of a heavy resonance, and/or through nonresonant production of multiple particles.

E. Effect of Finite $\lambda$

Still we have not captured all the physics. If $\lambda$ is infinite, there are no states in the spectrum with spin higher than two. But any confining gauge theory has hadrons with arbitrary spin. In gauge/string duality it has been found [23] that the ratio of the masses of hadrons with spin > 2 and higher to the lightest hadrons of spin 0 through 2 is $(\lambda)^{1/4}$. Thus there are towers of states of spin 5/2 and above whose lowest states lie near $(\lambda)^{1/4} \Lambda$. Note this need not be that large; if $N = 100$ one cannot take $\lambda > 100$ (see [61] for more details) and the fourth root of 100 is only about 3. Thus it is only legitimate to fully neglect these states, in this case, for processes which have center-of-mass energy below, say, 5$\Lambda$, or perhaps 10$\Lambda$. At higher energy, the higher-spin states are present, increasing the phase space for heavy resonance decays and the rate of multi-hadron production at large $q^2$. This drives us still further from the model of a narrow tower.

In addition, even a low-spin tower itself is altered by mixing of the original states in the tower with low-spin strings when the excitation level $n \sim \lambda^{1/4}$. This is a key breakdown of the gravity approximation to the string theory. Although the mixing is small, the number of strings at any given excitation level grows exponentially, so at large $n$ these states cannot be neglected. Unparticle production is then of massive strings in five dimensions, not of low-mass scalars in five dimensions. These massive strings then decay to a number of light resonances that lie at the bottom of the various towers.

This is almost enough to recover our one missing piece: the parton shower, and ensuing hadronization. This is claimed in [1] as the dominant high-energy process in a hidden valley, and is certainly the dominant process in QCD. Where is it in this five-dimensional language? More details will be given in [61], and the argument uses some subtle features of string theory, but I will simply claim here that the parton shower at large $\lambda$ is dual to a string falling in five dimensions toward the hard wall.

What is happening from the five dimensional point of view is easy to understand by looking carefully at what happens in the gauge theory. Let us use the language of QCD, speaking of color, quarks, antiquarks and gluons; we will then carry over the insight into the hidden sector. When a quark and antiquark are produced, they emit gluons, whose color lines are correlated with the quarks and with each other, as captured in 't Hooft’s double-line notation. As is well known, at any moment in time we can draw a line from the quark to the gluon whose anticolor is correlated with the quark’s color, and from there to the next gluon whose anticolor is correlated with the first gluon’s color, and so forth — forming a string. This is the same string which is used at the moment of fragmentation in the Lund string model, but before fragmentation occurs, during the parton shower, this is a string falling in five dimensions. (Note this is no classical string; it is a quantum, fluctuating, string.) It falls from $r \sim \sqrt{x/R^2}$ down to $r = \Lambda/R^2 = r_{\text{IR}}$, where it stops falling and hadronization occurs. As the string falls, an observer at a fixed $r$, corresponding to a probe of the string at some momentum scale $\propto r$, can resolve smaller and smaller structures within the string [58]. Thus the string contains more and more resolveable partons as the scale decreases.

Is this claim, that the full string theory is needed to see the parton shower, plausible? As another line of argument, which is really the same argument in a crossed channel, consider the following. The same operators which control parton splitting in DGLAP evolution control the leading-order behavior of parton splitting in the parton shower. We must keep the classically-twist-two
operators of spin 4 and higher at weak coupling, if we are to see the parton shower. If we then dial up the coupling, and drop all operators of dimension $\sim \lambda^{1/4}$, which include all operators of spin $>2$, we will be throwing away the parton shower itself. A more precise version of this argument will be presented elsewhere.

Of course, when the string reaches the bottom at $r_{IR} = \Lambda R^2$, it may for an instant be viewed as a highly excited, highly tangled hadron. It cannot be viewed as one of the hadrons in the tower of a simple five-dimensional scalar; such hadrons are five-dimensional points and have no internal structure. Here we have an extended object, with many quantum numbers, fundamentally a string. Interactions then cause the string to fragment, slowly if $N$ is large, quickly if $N$ is smaller or if the number of light flavors $N_f$ is nonzero. The hadronization process is, as in the Lund string model, the breaking apart of a string which is largely four-dimensional, being localized in the fifth dimension near the minimum value of $r_{IR}$.

In short, we must remember that $\lambda$ is finite, and include the strings, if we are to see the physics of parton showers and the process of hadronization. This physics can play an important role in the phenomenology, as outlined in [1].

\section{Conclusions}

In this section a more physical version of the (formerly)-narrow-tower model has been reconstructed. When modes of the hidden sector are produced, we expect the following signatures:

- At low energy $\sim \Lambda$, one may find a few light stable hidden hadrons of low spin, which decay back to standard model particles with long lifetimes.

- At higher energies one finds excited resonances, which decay rapidly to two or three light stable hidden hadrons, which in turn decay to standard model particles in the final state.

- Next the resonances become a continuum; production of several hidden hadrons becomes likely, leading to an increasing multiplicity of light hidden hadrons and consequently of standard model particles.

- At still larger energy, the parton shower begins to play a role in the evolution of the final state, making a purely hadronic description of the process inconvenient and indeed misleading, but without changing the basic signature of a high-multiplicity event.

- In all of these processes, the final state consists of a number of light neutral long-lived hidden-sector hadrons, with a multiplicity that grows with energy.

- The lightest hadrons have lifetimes orders of magnitude longer than most other hadrons, and may produce observable displaced vertices or give missing-energy signals.

The specific signals observed will of course depend on the details of its gauge and matter content of the hidden sector. But having included five-dimensional interactions and five-dimensional strings, we now see that this model has all the features and signatures of a confining hidden valley. As we did not rely upon strict conformal invariance, the result is largely independent of whether there are true unparticles at energies above the confinement scale.

However, despite the universality of the result, there are some important features of the hidden-valley signal which can be affected by strong dynamics. Let us now turn our attention to these.

\section{Other Effects of Strong Dynamics on Hidden Valley Phenomenology}

From the way I have presented things, the reader might be left with the impression that the ultraviolet strong dynamics that is present in conformal field theories with large anomalous dimensions has no impact whatsoever on the infrared physics. But this is not the case. Even if the low-lying states in the hidden sector are visible, giving hidden valley signals and making unparticle methods less central to the phenomenology, the strong dynamics can have a crucial impact on what we will see. Yet this comes not from the produced “unparticle”, but from other effects of strong coupling on the phenomenology: on resonances, on flavor symmetries within the hidden sector, on supersymmetry breaking within the hidden sector, and on the hidden parton shower.

\subsection{Effect of Strong Dynamics on Resonances}

As we noted in Sec. [I], one may well expect narrow resonances just below the point where continuum production of the hidden sector begins. (There may also be resonances, narrow or wide, within the continuum.) The spacing between the resonances, and the number of resonances that precede the continuum, are an important opportunity for learning about the hidden sector, just as charmonium was an important probe of QCD. Had QCD been more strongly coupled, the spacing between the charmonium states would have been wider, and the number of states below the open-charm continuum might have been very large. (Indeed in the large $\lambda$ limit, using [15], it was found [16] that a number of extremely deeply bound states may in some cases appear well below the onset of a continuum. The known examples require supersymmetric cancellations, but perhaps there are other mechanisms to obtain similar effects.)

In fact, if the resonances are from a massive field $\phi$ whose mass does not strongly destabilize the conformal
dynamics, then the $\phi$-onium states will still be bound by exchange of effectively-massless gauge bosons. This can happen in the toy models B and C of Sec. 11 if only one of the $N_f \gg 1$ scalars becomes massive at $M$, the others remaining light. The $\phi$-onium states will form an almost perfectly positronium-like system (though possibly at much stronger coupling) and the strong coupling can be measured from the positions of the resonances. Note this neither measures nor depends upon the dimension of any unparticle; conformal exchange is always Coulomb exchange, and only its coefficient depends on the coupling.

However, the central phenomenological question is whether and how the resonances can be observed. The branching fraction directly to two or three standard model particles will be low in this case, because annihilation to the effectively-massless hidden sector gauge bosons will dominate; see [4] for a brief discussion of hidden quarkonium. If the low-energy theory has some of the signatures of Sec. II E or II F, then the decays will be observable. Still, detailed kinematic reconstruction will have poor resolution at a hadron collider, so it is unclear how much information can be extracted without an ILC.

B. Effect of Strong Dynamics on Global Symmetries

In models with conformal dynamics and approximate global symmetries, renormalization effects can often either enhance or destroy those symmetries. This is determined by whether global-symmetry-breaking operators have larger or smaller dimensions than global-symmetry-preserving operators; if they are larger, then the global symmetry is accidentally restored at low energy, while if smaller then flavor is badly broken. For instance, in models A and B, a critical question is whether the adjoint or singlet scalar-bilinear has a larger anomalous dimension.

As an illustration, let us imagine how different the standard model might have been had it been conformal at high energy. Imagine for example that the standard model itself became strongly interacting above the electroweak scale. This is not a completely consistent example, but still instructive. A very important effect is that the Yukawa operators coupling the Higgs boson to standard model fermions would have nontrivial anomalous dimensions. If the dimensions of the Yukawa operators were all less than 4, than all Yukawa couplings would be driven large, all masses would be driven toward 100 GeV. But if all the operators were irrelevant (as in the simplest versions of technicolor) then all masses would end up small (which is in part why simple technicolor has trouble with the top quark Yukawa coupling.) If all masses are driven very small then there can be greatly enhanced symmetries: for example, light $c$, $b$ and $t$ quarks would give larger chiral symmetries below the QCD scale. This would then introduce stronger GIM-like suppressions into flavor-changing processes, and organize the hadrons into larger multiplets.

Another more subtle effect is that mixing angles between quarks could also be reduced by strong dynamics. If some Yukawa operators have very different anomalous dimensions from others, then the matrix of Yukawa couplings may tend to become highly structured and the CKM matrix is driven diagonal. Indeed, an analogous phenomenon was put to use in a realistic model of flavor in [52].

In a similar way, strong dynamics in the hidden sector can easily drive physics into regimes where hidden flavor symmetries are naturally approximate, rather than being either violated at order one or exact. Such a structure will in turn can have observable effects, such as

- A large multiplet of nearly degenerate hidden states which cannot decay to one another, and thus must decay via emission of standard model particles;
- Increased lifetimes for decays between hidden-sector “generations,” due to reduced intergenerational mixing angles.

Both of these can have a major impact on the phenomena observed.

Unfortunately, these signatures are not unique to, nor are they required by, strong dynamics, so they are not smoking guns for large anomalous dimensions. Also, I know of no new urgent experimental issues raised by this possibility that are not already under discussion. In short, though interesting, this consequence of strong coupling is something that theorists might consider further, but appears not to be urgent for experimentalists preparing for the LHC.

C. Effect of Strong Dynamics on Supersymmetry Breaking

It is well known that strong dynamics in a supersymmetric theory can suppress part or all of supersymmetry breaking in that sector; for an application to phenomenology, see [54] (which contains a review of the basic physics in an appendix.) Thus although supersymmetry breaking in our sector may have a characteristic scale of 100s of GeV, the hidden sector may be much more supersymmetric than our own. The spectrum of hidden sector particles is not easy to determine, but might show approximate degeneracies amongst bosons and fermions. Indeed it is even possible that the discovery of supersymmetric particles, or convincing verification of supersymmetry in nature, will occur through the hidden valley phenomenology.

For example, if in addition to hidden-hadron decays to the standard model there are supersymmetric hidden-hadrons (analogous to R-hadrons within QCD [54]), then the decay patterns of the various states may reveal degeneracies, and perhaps relations among decay rates or
branching fractions, that would be characteristic of a supersymmetric theory. One amusing example is that of $\phi$-onium resonances within a still-conformal sector, a case mentioned in Sec. IV A. There one could imagine predicting, and observing, effects from the supersymmetric generalization of positronium.

Let us make two immediate cautionary remarks. This scenario is not unique to strong dynamics. Supersymmetry breaking could be suppressed in the hidden sector by other, purely perturbative means. For example, in gauge mediation the messenger sector coupling to the hidden sector may be absent, or the relevant messengers very heavy, such that the standard model sector receives a louder message and gets a larger array of soft masses than does the hidden sector. Also, very weak gauge couplings in the hidden sector would result in rather weak supersymmetry-breaking effects. Furthermore, even if supersymmetry-breaking effects are large for most particles, approximate symmetries or dimensional counting arguments can easily make gaugino masses very small, leaving a hidden sector whose low energy physics might accidentally be an almost exactly supersymmetric Yang-Mills theory, with massless hidden gluons and with hidden gluinos light compared to the hidden confinement scale.

The other problem is that degeneracies in the hidden sector spectrum can arise from ordinary bosonic global symmetries, or simply from kinematics. For example, in a hidden pure Yang-Mills theory, the masses of metastable hidden glueballs which can only decay to the standard model sector lie within a narrow band, just by kinematics. In a model with light pions, the observable pions may have very similar masses because they transform as a simple multiplet under an accidental global symmetry — possibly enhanced through strong dynamics, as discussed above in Sec. IV B.

Thus, it is certainly possible that a hidden sector will reveal the first superpartners and other features of supersymmetry at the LHC, but it is likely to be far from obvious. Fortunately, it appears that there are no special experimental challenges in the supersymmetric limit, so that detection and analysis techniques do not need to be specially tuned. It appears to be enough to make careful but standard measurements of multiple processes.

D. Effect of Strong Dynamics on the Parton Shower

By contrast, the last effect I want to consider is of both theoretical and experimental importance. In hidden valley models with strong dynamics far above the mass gap, the parton shower can be very much more powerful than in QCD. This potentially can turn hard jets into soft spray, and make events more spherical, with much higher multiplicities, than in hidden valleys with weaker couplings above the mass gap. This possibility poses significant experimental challenges, whose details cannot be known without further theoretical development.

Let us first see why the parton shower is the most sensitive ingredient in a hidden valley to strong conformal or near-conformal dynamics in the 1–1000 GeV range. The first element in a hidden valley model is the coupling between the two sectors via some communicator, can be impacted by the strong dynamics, but indirectly, in a way that might not readily be observed. (There could be a large impact on the line shape of an easily observed particle, however.) The third element, a mass gap, explicitly involves the breaking of conformal invariance, and we have seen how the phenomenology can emerge in many different ways that are not especially sensitive to the dimensions of operators in the conformal regime. What of the second ingredient, the multi-particle production mechanism? Cascade decays explicitly involve breaking of conformal symmetry, and hadronization is a violent violation of conformal symmetry. But the parton shower is conformal dynamics in action.

The very form of the parton shower is determined by anomalous dimensions of operators, as noted above. These are not the low-spin low-dimension operators which we might couple to the standard model in the Lagrangian, but Wilson lines, or in the crossed channel, the high-spin high-dimension classical-twist-two operators which are always present in a gauge theory.

If these operators have small anomalous dimensions, the parton shower is inefficient. In QCD the quark in a $qq$ production process loses only a moderate fraction of its energy to gluon emission, much of which is collinear with the quark, maintaining a coherent jet. For this reason, $qq$ production gives predominantly two-jet events, sometimes three. In a weakly-coupled hidden valley, this jetty structure in the hidden sector is typically retained, though blurred, as the hidden hadrons decay into standard model particles. An example is shown in Fig. 25.

However, if the DGLAP operators have large anomalous dimensions, the parton shower is very efficient. (See 42 for a related effect in deep inelastic scattering.) The quarks and gluons shower so quickly, through both collinear and soft emission, that they all become soft. Soft emission is not collinear with the initial quark, and so the original direction of the quark’s motion is largely forgotten. Moreover, hard emission is also not supressed as it is in weak coupling, so the production process itself is likely to have several gluons along with the quark and antiquark. Altogether, the events at strong coupling are likely to be more spherical than jetty, though I do not know how to calculate the fluctuations away from perfect sphericity. 65 Also, the number of hidden hadrons will be larger, and their $p_T$ distribution much softer than at weak coupling. A guess at the appearance of such an event is shown in Fig. 26; again I emphasize this is a guess. 66 Notice how the preferred axis and the strongly variable calorimeter signal in Fig. 25 are absent in Fig. 26.
FIG. 25: An event (generated with HVMC 0.5) in which a 3.2 TeV $Z'$ decays to 30 GeV $\nu$-pions (see [1] for definitions) in a hidden sector which has a weak coupling above $\sim 100$ GeV. Notice the thrust axis is roughly vertical, though the events are by no means not pencil-like. The event shown contains roughly twenty bottom quarks and tau leptons.

FIG. 26: An event (generated with HVMC 0.5) in which a 3.2 TeV $Z'$ decays to 30 GeV $\nu$-pions (see [1] for definitions) in a hidden sector which has a strong coupling at all energies. Notice the event is now spherical. The event shown contains roughly fifty bottom quarks and tau leptons.

There are many challenges here which need to be addressed. For example, what will result from the application of jet algorithms to these events? How much energy must such events have such that they are easy to see? Are there any serious detector backgrounds which could mimic such an effect? What if most of particles produced are long-lived; could the pattern recognition software become too confused to operate properly? And if such events are identified, what questions can we ask of them to identify their source? What observables will most usefully allow analysis of such events? And finally, what is needed in formal theoretical development so that the guesswork that goes into Fig. [26] can be replaced by reliable prediction?

V. CONCLUSION

We have seen that unparticle models [5, 6, 31] with mass gaps [13, 14, 15] are typically examples of hidden valley models [1, 2, 3]. Not all hidden valley models are conformal; not all unparticle models have mass gaps which result in particles with observable decays; but there are many “hidden valley/unparticle” models with both features. We have seen that in these cases the dominant exclusive processes typically involve classic hidden valley phenomenology: new neutral light particles with long lifetimes, often produced with high multiplicity, along with new Higgs decays (and also decays of LSPs and related particles, though we did not study these here) to the hidden-valley particles. Some of the hidden-valley particles themselves decay to standard model particles, possibly with observably displaced vertices.

To see this required some tweaking of results in the literature, in particular, clarifying what can and cannot be said about conformal symmetry breaking using the language of unparticles [13], and adjusting the narrow-tower model of unparticles [15] to include various corrections which though small make a huge change in the physical phenomena. We also saw that the assumption that the hidden sector physics is invisible, as for example in [14], is often too pessimistic; the Higgs decays may not only be visible, they may be spectacular, as in [1, 2].

Still, in addition to this, inclusive studies using the above-mentioned events, and certain rare processes, may be able to reveal the special conformal kinematics associated to unparticles. I have also suggested that a dominant effect of strong coupling, and in particular large anomalous dimensions for twist-two operators, is the strong enhancement of the parton shower, and an increased sphericity, higher multiplicity, and lower $p_T$ spectrum in high-energy hidden valley events. More work is needed to confirm this suggestion, and to understand how strong dynamics, with or without conformal invariance, can affect the phenomenology of hidden valleys in other observable ways.

As in QCD, where both exclusive and inclusive questions have their merits, and where approximate scale invariance plays an important role in many analyses of QCD data, the study of a hidden sector should proceed by combining information from both exclusive final states (which, if visible, involve hidden valley phenomenology) and inclusive final states (which are determined by unparticle computations in the scale-invariant region, but
not elsewhere.) It may happen that all exclusive phenomena are invisible, and then one can only discuss the physics from the inclusive point of view. But if the exclusive events can be observed, they are typically more abundant, and are often more spectacular, more easily separated from background, and more informative. They are also often very unusual, and in some cases may pose serious challenges for the Tevatron and LHC experiments. These challenges should be addressed (and in some contexts are already being addressed) in the immediate future.

In conclusion, conformal invariance and inclusive signatures are powerful tools, but they are not powerful enough to address the physics of conformal-symmetry breaking, where the full diversity and complexity of quantum field theory may be found. Two-point functions of local composite operators, and the constraints of conformal invariance, simply cannot capture the phenomenological richness so often found in a hidden valley.

I am pleased to thank A. De Roeck, J.L. Feng, H. Lubatti, M. Graesser, B. Mele, S. Thomas, D. Ventura and K. Zurek for useful discussions.

[1] M. J. Strassler and K. M. Zurek, Phys. Lett. B 651, 374 (2007) [arXiv:hep-ph/0604261].
[2] M. J. Strassler and K. M. Zurek, arXiv:hep-ph/0605193.
[3] M. J. Strassler, arXiv:hep-ph/0607160.
[4] T. Han, Z. Si, K. M. Zurek and M. J. Strassler, arXiv:0712.2041 [hep-ph].
[5] H. Georgi, Phys. Rev. Lett. 98, 221601 (2007) [arXiv:hep-ph/0703260].
[6] H. Georgi, Phys. Lett. B 650, 275 (2007) [arXiv:0704.2457 [hep-ph]].
[7] N. Seiberg, Nucl. Phys. B 435, 129 (1995) [arXiv:hep-th/9411149].
[8] L. Randall and R. Sundrum, Phys. Rev. Lett. 83, 4690 (1999) [arXiv:hep-th/9906064].
[9] L. Randall and R. Sundrum, Phys. Rev. Lett. 83, 3370 (1999) [arXiv:hep-ph/9905221].
[10] J. J. van der Bij and S. Dilcher, Phys. Lett. B 638, 234 (2006) [arXiv:hep-ph/0605008].
[11] Z. Chacko, H. S. Goh and R. Harnik, Phys. Rev. Lett. 96, 231802 (2006) [arXiv:hep-ph/0506256].
[12] G. Burdman, Z. Chacko, H. S. Goh and R. Harnik, JHEP 0702, 009 (2007) [arXiv:hep-ph/0601512].
[13] P. J. Fox, A. Rajaraman and Y. Shirman, Phys. Rev. D 76, 075004 (2007) [arXiv:0705.3092 [hep-ph]].
[14] A. Delgado, J. R. Espinosa and M. Quiros, JHEP 0710, 004 (2007) [arXiv:0707.4309 [hep-ph]].
[15] M. A. Stephanov, Phys. Rev. D 76, 035008 (2007) [arXiv:0705.3049 [hep-ph]].
[16] S. Chang, P. J. Fox and N. Weiner, JHEP 0608, 068 (2006) [arXiv:hep-ph/0511250].
[17] S. Chang, P. J. Fox and N. Weiner, Phys. Rev. Lett. 98, 111802 (2007) [arXiv:hep-ph/0608310].
[18] J. R. Espinosa and J. F. Gunion, Phys. Rev. Lett. 82, 1084 (1999) [arXiv:hep-ph/9807275].
[19] U. Ellwanger, J. F. Gunion, C. Hugonie and S. Moretti, arXiv:hep-ph/0305109, [arXiv:hep-ph/0401228]. R. Dermisek and J. F. Gunion, Phys. Rev. Lett. 95, 041801 (2005) [arXiv:hep-ph/0502105]. U. Ellwanger, J. F. Gunion and C. Hugonie, JHEP 0507, 041 (2005) [arXiv:hep-ph/0503203]. R. Dermisek and J. F. Gunion, arXiv:hep-ph/0510322.
[20] L. M. Carpenter, D. E. Kaplan and E. J. Rhee, arXiv:hep-ph/0607294. C. Morningstar and M. J. Peardon, Nucl. Phys. Proc. Suppl. 83, 887 (2000) [arXiv:hep-lat/9911003].
[21] C. J. Morningstar and M. J. Peardon, Phys. Rev. D 60, 034509 (1999) [arXiv:hep-lat/9901004]. C. Morningstar and M. J. Peardon, Nucl. Phys. Proc. Suppl. 83, 887 (2000) [arXiv:hep-lat/9911003].
Rev. D 74, 015005 (2006) [arXiv:hep-ph/0602229].

[44] R. C. Brower, J. Polchinski, M. J. Strassler and C. I. Tan, arXiv:hep-th/0603115.

[45] A. Karch and E. Katz, JHEP 0206, 043 (2002) [arXiv:hep-th/0205236].

[46] M. Kruczenski, D. Mateos, R. C. Myers and D. J. Winters, JHEP 0307, 049 (2003) [arXiv:hep-th/0304032].

[47] J. L. Hovdebo, M. Kruczenski, D. Mateos, R. C. Myers and D. J. Winters, Int. J. Mod. Phys. A 20, 3428 (2005).

[48] M. Bander, J. L. Feng, A. Rajaraman and Y. Shirman, Phys. Rev. D 76, 115002 (2007) [arXiv:0706.2677 [hep-ph]].

[49] S. Hong, S. Yoon and M. J. Strassler, JHEP 0604, 003 (2006) [arXiv:hep-th/0409118].

[50] G. F. de Terramond and S. J. Brodsky, Phys. Rev. Lett. 94, 201601 (2005) [arXiv:hep-th/0501022].

[51] S. Hong, S. Yoon and M. J. Strassler, arXiv:hep-ph/0501197.

[52] R. C. Brower, J. Polchinski, M. J. Strassler and C. I. Tan, arXiv:hep-th/0603115.

[53] J. Erlich, E. Katz, D. T. Son and M. A. Stephanov, Phys. Rev. Lett. 95, 261602 (2005) [arXiv:hep-ph/0501128].

[54] A. E. Nelson and M. J. Strassler, JHEP 0207, 021 (2002) [arXiv:hep-ph/0104051].

[55] A. E. Nelson and M. J. Strassler, JHEP 0009, 030 (2000) [arXiv:hep-ph/0006251].

[56] G. R. Farrar, arXiv:hep-ph/9408379.

[57] S. Gopalakrishna, S. Jung and J. D. Wells, in preparation.

[58] L. Susskind, Phys. Rev. D 49, 6606 (1994) [arXiv:hep-th/9308139].

[59] D. T. Son and M. A. Stephanov, Phys. Rev. D 69, 065020 (2004) [arXiv:hep-ph/0304182].

[60] N. Arkani-Hamed, A. G. Cohen and H. Georgi, Phys. Rev. Lett. 86, 4757 (2001) [arXiv:hep-th/0104005].

[61] M.J. Strassler, unpublished.

[62] M.J. Strassler, unpublished.

[63] Deconstruction [60] refers to the discretizing of a space, as in [59], and representing it as a gauge theory in one lower dimension. The number of Kaluza-Klein modes in a deconstructed theory is finite. The introduction of an infrared cutoff while retaining the continuity of the five-dimensional space is a form of compactification; the number of Kaluza-Klein modes is countably infinite. In this context this cutoff is known as the “hard wall” model, a model often used in the description of a confining gauge theory, which is why “Randall-Sundrum 1” [9] can be viewed as dual to technicolor.

[64] This event was generated using the HVMC Monte Carlo, version 0.5 [62]. This Monte Carlo is based on Pythia 6.4, combining its routines to simulate $Z'$ decay to the hidden sector, showering and hadronization within the hidden sector, and decay of hidden-sector hadrons back to standard model particles, which is followed by standard model showering and hadronization.

[65] The statements made here are clearest at large $N$; there will be $1/N^2$ corrections to these statements that deserve more study, in which color singlet combinations of partons are radiated off to begin their own, separate parton shower. Whether these could introduce strong fluctuations in the appearance of the events is not yet clear.

[66] A $Z'$ decay to a hidden quark and antiquark was dressed with a Pythia parton shower of hidden gluons, in which the showering rate was enhanced by fixing the gauge coupling at a large enough value that collinear effects largely vanished by the scale of hidden confinement. The result may bear only a passing resemblance to the actual physics of a strongly coupled theory, but this guess appears physically reasonable, and should at least be thought-provoking.