PHYSICAL REVIEW D **76,** 035008 (2007)

Deconstruction of unparticles

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(Received 19 June 2007; published 31 August 2007)

We discuss properties of hypothetical scale invariant (unparticle) matter by viewing it as a tower of massive particles. We show how peculiar properties of unparticles emerge in the limit when the mass spacing parameter Δ vanishes. We explain why the unparticle cannot decay in this limit and how, for finite Δ, the decays manifest themselves in a relation between the reconstructed invariant mass and vertex displacement. We describe a model field theory in AdS_5 which explicitly implements the deconstruction procedure by truncating the extra dimension to size of order $1/\Delta$.

DOI: 10.1103/PhysRevD.76.035008 PACS numbers: 14.80.−j, 11.15.Tk

I. INTRODUCTION

In a recent paper [1] Georgi suggests the consideration of a hypothetical scale invariant, or conformal, matter very weakly coupled to the standard model matter as a possible component of physics above TeV scale which might begin to show up at the LHC. A vast amount of knowledge exists about the properties of conformal field theories, and its applications to condensed matter problems, such as critical phenomena, abound. However, there is little understanding of how such a conformal sector, if it exists, would manifest itself in particle physics experiments. Georgi terms such matter "unparticle," for its properties and signatures are qualitatively different from those of particles. Several curious properties of unparticles have been exposed in Ref. [1,2], including the unusual scaling of the apparent phase space volume, the unusual missing energy spectra, unusual interference patterns, etc. Several recent papers have been addressing novel signatures of unparticles [3–9].

The purpose of this paper is to clarify the notion of the unparticle using the language familiar to a particle physicist. To that end we deconstruct the unparticle and view it as an infinite tower of particles of different masses.[1] We show how the peculiar properties of unparticles exposed in Refs. [1,2] can be interpreted and understood in that language.

We shall think of the unparticle as a limiting case in which the spacing Δ^2 of the (squared) masses in the tower of particles goes to zero. Using this limiting procedure we explain an apparent paradox mentioned in Ref. [2], that the unparticle need not decay despite the presence of a finite imaginary part in its two-point correlator. As a corollary we find that if Δ is small but finite, the unparticle can decay, and we describe peculiar signatures of such decays.

In the beginning we shall view deconstruction as a purely mathematical device translating the properties of the unparticle into particle physics language. In Sec. VI we shall take a more constructive approach and describe a model of a field theory in which the unparticle indeed arises as a limiting procedure. Not surprisingly, the model requires an additional space dimension, and the tower of deconstructing particles appears naturally as a Kaluza-Klein tower, once the extra dimension is compactified/truncated. In order to produce an unparticle with noninteger scaling dimension we use AdS_5 geometry and describe the unparticle using a massive scalar field.

II. SETUP AND NOTATIONS

Following Ref. [1], let us imagine that there exists a scale invariant sector of our world described, e.g., by some strongly self-coupled conformal theory. Scale invariance means that there are no (massive) particles in this sector. Since such an unparticle sector is not seen in present experiments, we must further assume that the coupling of this conformal sector to the standard model particles/fields is very weak—somewhat in the spirit of the "hidden-valley" models [10]. This interaction can be described in an effective field theory language as a (nonrenormalizable) coupling between some standard model operator and the unparticle operator O of scaling dimension $d_{\mathcal{U}}$. Since the unparticle sector is self-interacting, the dimension of O can be nontrival (noninteger).

The correlation function of the operator is given by

$$\int d^4x e^{iPx}\langle 0|TO(x)O^\dagger(0)|0\rangle = \int \frac{dM^2}{2\pi}\rho_O(M^2) \times \frac{i}{P^2 - M^2 + i\varepsilon}. \quad (1)$$

By scale invariance the spectral function of the operator O must be a power of M^2:

$$\rho_O(M^2) = A_{d_{\mathcal{U}}}(M^2)^{d_{\mathcal{U}}-2}, \quad (2)$$

where $A_{d_{\mathcal{U}}}$ is a normalization constant chosen by convention in Ref. [1]. Its precise form/value is not consequential

[1] We rely on a more generic meaning of the term "deconstruction," and wish to emphasize the difference from the so far more common specific use of this term to describe deconstruction of five-dimensional theories by discretization of the fifth dimension.

for the discussion below. On the other hand,

$$\rho_O(M^2) = 2\pi \sum_\lambda \delta(M^2 - M_\lambda^2)|\langle 0|O(0)|\lambda\rangle|^2. \quad (3)$$

The sum in Eq. (3) is over all relativistically normalized states $|\lambda\rangle$ (at fixed spatial momentum—i.e., no $d^3 p_\lambda$ integration). The unparticle spectral function means that the spectrum of the operator O is continuous, i.e., the sum in Eq. (3) is in fact an integral. Let us imagine that the scale invariance is broken in the system in a controllable way, so that, instead of a continuous spectrum of states λ, there is a discrete tower of states with the spacing controlled by parameter Δ. To simplify discussion, we shall pick a particular spectrum

$$M_n^2 = \Delta^2 n. \quad (4)$$

It is straightforward to adjust all subsequent discussion to any other spectrum, e.g., $M_n^2 = \Delta^2 n^2$, etc. We shall assume that Δ is much smaller than other scales pertinent to the problem. Let us introduce notation for the matrix element for the nth particle:

$$F_n^2 \equiv |\langle 0|O(0)|\lambda_n\rangle|^2. \quad (5)$$

We can then write:

$$\rho_O(M^2) = 2\pi \sum_n \delta(M^2 - M_n^2) F_n^2 \quad (6)$$

and

$$\int d^4x e^{iPx}\langle 0|TO(x)O^\dagger(0)|0\rangle = \sum_n \frac{iF_n^2}{P^2 - M_n^2 + i\varepsilon}. \quad (7)$$

In the limit $\Delta \to 0$ the sum over n in Eq. (6) becomes an integral, which must match Eq. (2). From this condition we easily determine that F_n must be given by

$$F_n^2 = \frac{A_{d_\mathcal{U}}}{2\pi}\Delta^2 (M_n^2)^{d_\mathcal{U}-2}. \quad (8)$$

The constants F_n are similar to the decay constants of mesons in QCD. More generally,

$$M_n^2 = \Delta^2 n^{1/\gamma} \Rightarrow F_n^2 = \frac{A_{d_\mathcal{U}}}{2\pi\gamma}\Delta^{2\gamma}(M_n^2)^{d_\mathcal{U}-2-(\gamma-1)}. \quad (9)$$

III. PRODUCTION

Consider the example from Ref. [1] of production of such an unparticle. Imagine the coupling of the unparticle given by

$$i\frac{\lambda}{\Lambda_\mathcal{U}^{d_\mathcal{U}}}\bar{u}\gamma_\mu(1-\gamma_5)t\partial^\mu O + \text{H.c.}, \quad (10)$$

where λ is a dimensionless coupling and $\Lambda_\mathcal{U}$ is the Banks-Zaks scale in the unparticle theory. Using representation (7) we can define the deconstructing particle field

$$\lambda_n(x) \equiv O(x)/F_n. \quad (11)$$

According to (7), on the mass shell of the n's particle this field will be canonically normalized. Thus, the interaction Eq. (10) becomes after deconstruction

$$i\frac{\lambda}{\Lambda_\mathcal{U}^{d_\mathcal{U}}}\bar{u}\gamma_\mu(1-\gamma_5)t\sum_n F_n\partial^\mu\lambda_n + \text{H.c.} \quad (12)$$

Now it is easy to study the production of the unparticle using the standard notions of particle physics. The kinematics is that of a two-body decay of a t quark. That is, for each n the energy of the u quark is fixed to $E_u = (m_t^2 - M_n^2)/(2m_t)$. The spectrum of E_u consists of a peak for each value of n, which in the limit $\Delta \to \infty$ merge into the continuum distribution displayed in Ref. [1]. The decay rate for each n is

$$\Gamma(t \to u + \lambda_n) = \frac{|\lambda|^2}{\Lambda_\mathcal{U}^{2d_\mathcal{U}}}\frac{m_t E_u^2}{2\pi}F_n^2. \quad (13)$$

An interval dE_u corresponds to an interval of masses $dM^2 = 2m_t dE$ which contains $2m_t dE/\Delta^2$ states λ_n. Thus we obtain

$$\frac{d\Gamma}{dE_u} = \frac{2m_t}{\Delta^2}\Gamma(t \to u + \lambda_n) = \frac{2m_t}{\Delta^2}\frac{|\lambda|^2}{\Lambda_\mathcal{U}^{2d_\mathcal{U}}}\frac{m_t E_u^2}{2\pi}\left(\frac{A_{d_\mathcal{U}}}{2\pi}\Delta^2(M_n^2)^{d_\mathcal{U}-2}\right), \quad (14)$$

with $M_n^2 = m_t^2 - 2m_t E_u$, in agreement with Ref. [1]. We see that each of the deconstructing particles λ_n couples weaker and weaker as $\Delta \to 0$, but their number in a fixed interval of energies dE_u is increasing inversely proportionally to their coupling leading to finite $d\Gamma/dE_u$ in the scaling limit $\Delta \to 0$.

IV. INTERFERENCE

Another example considered in Ref. [2] is the coupling of a vector unparticle operator O^μ to a neutral vector or axial lepton current, e.g., ($\ell = e$ or μ)

$$c_{A\mathcal{U}} M_Z^{1-d_\mathcal{U}}\bar{\ell}\gamma^\mu\gamma_5\ell O_\mu, \quad (15)$$

where, following notations in Ref. [2] we expressed the dimensionful coupling in units of the Z boson mass M_Z. These couplings produce contributions to, e.g., $e^+e^- \to \mu^+\mu^-$ amplitudes due to the virtual unparticle which interferes with the standard model γ and Z boson amplitudes. In the case of the vector operator O^μ, which we assume to be conserved $\partial_\mu O^\mu = 0$, the Eqs. (1) and (2), etc. generalize as

$$\Pi^{\mu\nu}(q) \equiv \int d^4x e^{iqx}\langle 0|TO^\mu(x)O^\nu(0)|0\rangle = (-g^{\mu\nu} + q^\mu q^\nu/q^2)\int \frac{dM^2}{2\pi}\rho_O(M^2) \times \frac{i}{q^2 - M^2 + i\varepsilon}, \quad (16)$$

where again by scale invariance ρ_O is given by Eq. (2). Evaluating the integral over M^2 one finds

$$\Pi^{\mu\nu}(q) = (-g^{\mu\nu} + q^\mu q^\nu/q^2)\frac{iA_{d_\mathcal{U}}}{2\sin(\pi d_\mathcal{U})} \times (-q^2 - i\varepsilon)^{d_\mathcal{U}-2}. \quad (17)$$

Deconstructing the unparticle operator O^μ proceeds similarly to the scalar operator. We introduce decay constants F_n via

$$\langle 0|O^\mu(0)|\lambda_n\rangle = \epsilon^\mu F_n, \quad (18)$$

where ϵ^μ is the polarization of the massive vector particle λ_n. Then the correlation function is given by

$$\Pi^{\mu\nu}(q) = (-g^{\mu\nu} + q^\mu q^\nu/q^2)\sum_n \frac{iF_n^2}{q^2 - M_n^2 + i\varepsilon}. \quad (19)$$

If we assume the same mass spectrum as in Eq. (4), the constants F_n are again given by Eq. (8).

The contribution to the $e^+e^- \to \mu^+\mu^-$ amplitude from the unparticle is proportional to the correlation function (16), and following Ref. [2] we define:

$$\Delta_\mathcal{U} = \frac{A_{d_\mathcal{U}}}{2\sin(\pi d_\mathcal{U})}(-q^2 - i\varepsilon)^{d_\mathcal{U}-2}. \quad (20)$$

This amplitude interferes with the amplitude due to the virtual Z proportional to

$$\Delta_Z = \frac{1}{q^2 - M_Z^2 + iM_Z\Gamma_Z}. \quad (21)$$

This standard model amplitude is mostly real away from the Z pole and is mostly imaginary near the pole. The unusual property of the unparticle amplitude (20) pointed out in Refs. [2,3] is that it has a nonzero imaginary part for all $q^2 > 0$. This allows the amplitudes $\Delta_\mathcal{U}$ and Δ_Z to interfere even at the Z pole, where the latter is imaginary.

This property follows naturally from the deconstructed picture in which

$$\Delta_\mathcal{U} = \sum_n \frac{F_n^2}{q^2 - M_n^2 + i\varepsilon}. \quad (22)$$

The imaginary part of the amplitude $\Delta_\mathcal{U}$ as a function of q^2 consists of a series of δ-function peaks at $q^2 = M_n^2$:

$$\text{Im}\,\Delta_\mathcal{U} = -\sum_n F_n^2\pi\delta(q^2 - M_n^2). \quad (23)$$

Each peak becomes lower as $F_n^2 \sim \Delta^2 \to 0$, but their density increases. Converting the sum over n into the integral over M_n^2 we find that

$$\text{Im}\,\Delta_\mathcal{U} \to -\frac{F_n^2}{\Delta^2}\pi = -\frac{A_{d_\mathcal{U}}}{2}(M_n^2)^{d_\mathcal{U}-2} \quad (24)$$

in agreement with (20). The factor $\sin(\pi d_\mathcal{U})$ which cancels in (20) never appears in the first place in (24).

Away from the Z pole, where Δ_Z is real, the interference term is proportional to $\text{Re}\Delta_\mathcal{U}$. This is given by the sum in (22) where particles with masses $M_n^2 < q^2$ contribute with the opposite sign from those with $M_n^2 > q^2$. The case $d_\mathcal{U} = 3/2$ is special, as pointed out in Ref. [2]: $\text{Re}\Delta_\mathcal{U} \sim \cot(\pi d_\mathcal{U})$ vanishes. This has a simple meaning—at this value of $d_\mathcal{U}$ particles with M_n^2 above q^2 exactly cancel contribution of particles below q^2 (for any q^2). This is most clear from the integral representation:

$$\text{Re}\,\Delta_\mathcal{U} = \int_0^\infty dM^2\frac{(M^2)^{d_\mathcal{U}-2}}{q^2 - M^2}. \quad (25)$$

That this (principal value) integral vanishes at $d_\mathcal{U} = 3/2$ can be seen by doing the change of variables $M \to q^2/M$ (mass inversion) which maps the regions above and below q^2 onto each other.

V. DECAY?

We observe (8) that each deconstructing particle λ_n couples with strength proportional to $F_n^2 \sim \Delta^2$ which vanishes as $\Delta \to 0$. Thus, in a certain sense, a true ($\Delta = 0$) unparticle, once produced, never decays. This limiting procedure explains the apparent paradox pointed out in Ref. [2]: the finite imaginary part (24) of the "propagator" of the unparticle does not mean it has a finite lifetime.

What if the unparticle sector is almost conformal with a very small but nonzero Δ? The lifetime of a deconstructing particle λ_n is proportional to $F_n^{-2} \sim \Delta^{-2}$, and let us assume that it is in the range that one can observe the displaced vertex of λ_n decay into ordinary standard model particles. What would the signatures of such decays be? For simplicity, let us assume here no interference with standard model amplitudes.

First of all, the invariant mass spectrum of the decay products (e.g., lepton pairs) will not peak but will be a monotonous distribution (we assume that Δ is much less than the experimental resolution). Furthermore, the lifetime would be proportional to F_n^{-2}, which depends on M_n according to Eq. (8) or (9). There are of course trivial kinematic and coupling factors, which might add an integer power of M_n. One would therefore observe secondary vertices whose average displacement is correlated with the invariant mass of the products of decay.

For example, the contribution of the interaction (15) to the decay rate of λ_n is (taking F_n from Eq. (8))

$$\Gamma(\lambda_n \to \mu^+\mu^-) = \frac{c_A^2 \mathcal{U} M_Z^{2-2d_\mathcal{U}}}{8\pi} F_n^2 M_n = \frac{c_A^2 \mathcal{U} M_Z^{2-2d_\mathcal{U}} A_{d_\mathcal{U}}}{16\pi^2} \Delta^2 M_n^{2d_\mathcal{U}-3}. \quad (26)$$

Thus, the lifetime τ_d measured through the mean displacement of a vertex ($\tau_d = \ell_d/(\gamma v)$), if observed, will scale with the reconstructed invariant mass as

$$\tau_d = 1/\Gamma \sim M^{3-2d_\mathcal{U}}. \quad (27)$$

VI. MODELING AND DECONSTRUCTING THE UNPARTICLE USING AdS$_5$

So far we have viewed deconstruction as an abstract mathematical trick to cast the interaction of the unparticle as a sum of the interactions of the particles λ_n. This construction can be made more explicit by considering a model of the unparticle based on a five-dimensional field theory. The idea is simple: restricting the extent of the fifth dimension to size of order $1/\Delta$ will lead to the necessary discrete spectrum of deconstructing particles λ_n.

For concreteness and simplicity let us focus on a scalar unparticle as in Sec. III. The correlator (7) can be obtained from a two-point Green's function of a massive scalar field $\Phi(x, z)$, where x is a Minkowski coordinate, while the fifth coordinate z can be thought of either as a continuous index, or as a fifth coordinate. We shall take the anti-de Sitter (AdS) metric for this five-dimensional space: $ds^2 = (dx_\mu dx^\mu - dz^2)/z^2$. The Lagrangian (density in Minkowski space) reads:

$$\mathcal{L} = \int dz \sqrt{g}[g^{MN}\partial_M\Phi\partial_N\Phi - m_5^2\Phi^2]/2, \quad (28)$$

where, as usual, $x^M = (x^1, x^2, x^3, t, z)$ and $g = \det\|g_{MN}\|$, $g_{tt} = +1$. Note that the mass parameter m_5 is dimensionless, and so is the field Φ. The operator O can then be defined in terms of the field Φ as follows:

$$O(x) \equiv \lim_{z\to 0} z^{-d_\mathcal{U}}\Phi(x, z). \quad (29)$$

In other words, the standard model operators such as, e.g., $\partial_\mu(\bar{u}\gamma^\mu(1-\gamma_5)t)$ from Eq. (10) couple to the field $\Phi(x, z)$ only on the boundary $z \to 0$.

The dimension of the operator $d_\mathcal{U}$ determines the required mass of the field (or is determined if the mass is given) by the well-known formula:

$$m_5^2 = d_\mathcal{U}(d_\mathcal{U} - 4). \quad (30)$$

The following analysis of this model bears obvious resemblance to the holographic technique described in [11,12] and developed in many subsequent works. In fact, our five-dimensional model could be perceived as a dual description of some four-dimensional conformal field theory in the sense of the AdS/CFT correspondence [13]. There is a similarity with the extra-dimensional scenarios [14,15], but here we consider a *scalar* field in the bulk, rather than gravity. Gauge fields have been also extensively studied in AdS$_5$ (see, e.g., [16–18] and refs. therein), as well as massless scalar field [19–22] in a similar setup, but different contexts. Here we shall focus on the case of the *massive* scalar field which will allow a nontrivial scaling dimension $d_\mathcal{U}$. With the understanding that many elements of the following analysis can be found in the above literature, we shall, nevertheless, carry the following discussion in a self-contained manner.

To understand and derive the relationship (30) between the rescaling factor in Eq. (29) and the mass parameter given by (30), let us recall that the two-point correlation function of the field Φ which appears in

$$\langle O(x)O(0)\rangle = \lim_{z,z'\to 0} z^{-d_\mathcal{U}}(z')^{-d_\mathcal{U}}\langle\Phi(x, z)\Phi(0, z')\rangle \quad (31)$$

is the Green's function of the linear differential operator obtained by taking two variational derivatives of the Lagrangian (28) with respect to Φ:

$$[\partial_z z^{-3}\partial_z + z^{-3}q^2 - z^{-5}m_5^2]G(q; z, z') = \delta(z - z'), \quad (32)$$

where $G(q; z, z')$ is the Fourier transform of $\langle\Phi(x, z)\Phi(0, z')\rangle$ with respect to x. The behavior of the Green's function at small z is easy to find by noticing that the term q^2 is negligible for $z \ll q^{-1}$, and that the Eq. (32) with q^2 neglected is solved (for $z \neq z'$) by a power ansatz $G \sim z^\sigma$, with $\sigma(\sigma - 4) = m_5^2$. Thus, on the account of Eq. (30): $G(q; z, z') \sim z^{d_\mathcal{U}}(z')^{d_\mathcal{U}}$.[2] Thus the correlator $\langle O(x)O(0)\rangle$ in Eq. (31) is finite in the limit $z, z' \to 0$, and by dimension counting must be proportional to $x^{-2d_\mathcal{U}}$, which is what the dimension of the operator $O(x)$ implies.

Consider now AdS space with finite extent in the z direction (AdS slice): $z \in [0, z_m]$, with $z_m \sim 1/\Delta$. We can write the representation for the Green's function in terms of the orthogonal set of normal modes ϕ_n:

$$G(q; z, z') = \sum_n \frac{\phi_n(z)\phi_n(z')}{q^2 - M_n^2 + i\varepsilon}, \quad (33)$$

normalized according to

$$\int dz z^{-3}\phi_n^2 = 1. \quad (34)$$

The modes behave as $\phi_n \sim z^{d_\mathcal{U}}$ for small z, and the constants F_n can be identified, comparing (7), (31), and (33), as

$$F_n = \lim_{z\to 0} z^{-d_\mathcal{U}}\phi_n(z). \quad (35)$$

This result is similar to the expression for the meson decay constants in AdS/QCD in terms of the $z \to 0$ asymptotics of the normalizable modes [23,24].

[2]The explicit solution, which we do not need here, is $G(q; z, z') = (\pi/2)(zz')^2 J_{d_\mathcal{U}-2}(qz)Y_{d_\mathcal{U}-2}(qz)\theta(z'-z) + (z \leftrightarrow z')$.

The distribution of masses M_n can be obtained given the specific boundary conditions at z_m. Instead of relying on the exact solution, let us note that only $n \gg 1$ modes interest us here, since, by assumption, $\Delta \ll M_n$. The equation for normal modes (32) can be cast into Schrödinger form by substitution: $\phi_n = z^{3/2}\psi_n$, and then solved in the WKB approximation. For $z \gg M_n^{-1}$, including the $z = z_m$ boundary, the large n modes are $\psi_n \sim \sin(M_n z - C_1)$, where constant C_1 depends on m_5 (i.e., on $d_{\mathcal{U}}$), but not on n. Thus the mass spectrum is given by $M_n \sim (z_m)^{-1} \times (\pi n + C_2)$, where constant C_2 is related to C_1 and depends also on the type of the boundary condition at z_m (e.g., $C_2 = C_1$ for the Dirichlet boundary condition). Neglecting $\mathcal{O}(1/n)$ terms we find quite generally $M_n^2 \to \Delta^2 n^2$ for $n \gg 1$, where $\Delta = \pi/z_m$. To obtain the linear spectrum as in (4) one can modify the AdS background at large z, as it is done in Ref. [25], instead of cutting the space off at z_m.

The arguments in this section assume that the dimension of the operator satisfies $1 < d_{\mathcal{U}} < 2$. Indeed, for the normal modes to be normalizable in the sense of Eq. (34) we must have $d_{\mathcal{U}} > 1$. More subtly, for the Green's function in Eq. (32) to have the behavior $z^{d_{\mathcal{U}}}(z')^{d_{\mathcal{U}}}$ as $z, z' \to 0$, the value $d_{\mathcal{U}}$ must be the *smallest* of the two solutions of quadratic Eq. (30), which means $d_{\mathcal{U}} < 2$. The integer values $d_{\mathcal{U}} = 1, 2$ are, of course, special and, although interesting, will not be considered here.

It should be also pointed out that the condition $m_5^2 > 0$, violated for $1 < d_{\mathcal{U}} < 2$ according to Eq. (30), is not at all necessary for the stability of the scalar theory in AdS space [26,27]. One way to see this is to observe that the differential operator $[-\partial_z z^{-3}\partial_z + z^{-5}m_5^2]$ in Eq. (32) is positive definite even for negative $m_5^2 > -4$, i.e., there are no unstable modes.

In summary, we have seen how the notion of the unparticle can be somewhat demystified by representing it as an infinite tower of massive particles with controllable mass-squared spacing Δ^2. We used such a deconstruction technique to rederive and clarify the peculiar properties of the unparticle pointed out in Refs. [1,2] and to show that the pure ($\Delta = 0$) unparticle cannot decay, while for small but nonzero Δ the decay is possible, with a peculiar signature. Finally, we described a possible field theory realization of the deconstruction procedure using a slice of AdS_5 space.

ACKNOWLEDGMENTS

The author thanks Professor W. Y. Keung for bringing Ref. [1] to his attention and for stimulating discussions. This research is supported by the DOE Grant No. DE-FG0201ER41195.

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