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BEYOND THE STANDARD MODEL WITH STRONG DYNAMICS

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Abstract

Field theoretical extensions of the Standard Model which retain its successful structural features and simplicity are analyzed and characterized. A general classification is provided, in particular, for theories with new strong dynamics where the Dark Matter candidate is an accidentally stable bound state.

1 The Standard Model paradigm

The Standard Model (SM) of particle physics is able to explain a vast amount of experimental data with remarkable precision, yet it is believed to be an effective description of a more fundamental theory. There are two main experimental facts which justify and in fact impose this attitude: the existence of Dark Matter (DM), and the impossibility to explain Baryogenesis in the context of the Standard Model. The existence of a Landau pole for hypercharge, and possibly the inclusion of gravity at the quantum level, represent, on the other hand, internal obstructions to consider the SM as a complete theory valid at all scales. Finally, there are hints which are also suggestive of a deeper theoretical layer, though they do not point to any internal inconsistency: *i)* the fact that SM fermions fill GUT multiplets, together with approximate gauge coupling unification; and *ii)* the observation that we seem to live in a very special point of the SM parameter space, away from which our universe would not have the rich chemistry that we observe and nuclei, including the proton, would be unstable. [1)]

The success of the SM as an Effective Field Theory by and large follows from two ingredients. First, its conservation laws and selection rules (such as baryon and lepton number conservation and approximate flavor and custodial invariance) naturally arise as the consequence of accidental symmetries emerging in the infrared, provided the cutoff scale is sufficiently high. The same indication for a large cutoff scale also comes from gauge coupling unification. Second, fermions come in complex (i.e. chiral) representations of the gauge group, so that their masses are explained dynamically in terms of couplings. It is extremely suggestive that no fermion has been observed which could have a bare mass prior to symmetry breaking.

A known unsatisfactory feature of the SM is that the electroweak (EW) scale is not predicted but rather derived from an input parameter. Moreover, reproducing the experimentally observed value requires a delicate tuning in light of quantum corrections. Finding an explanation of the EW scale within a natural extension of the SM has driven most of the efforts of the high-energy community in the last forty years. As a matter of fact, all the theories of this kind constructed so far lack the simplicity of the SM. Although they are able –by construction– to explain naturally the weak scale, they require additional assumptions to agree with other experimental observations. It is possible, on the other hand, that naturalness is not the right principle to follow to explain the electroweak scale, and that other mechanisms (such as anthropics [2)], criticality [3)], or cosmological relaxation [4)]) are at work.

With this in mind, it is reasonable to ask if minimal extensions of the SM exist which retain its virtues and simplicity and where the stability of the Dark Matter candidate follows as the result of an accidental symmetry.[1] Taking seriously the hint of gauge couplings unification, one could also ask the new theory to explain it in terms of a grand unified phase at high energies. The prototype of such class of models is one where the SM is extended by new gauge dynamics and new fermion fields in complex representations ("dark" sector).[2] In the following we will focus on the case in which the dark gauge group (or part of it) gets strong in the infrared and confines, so that the DM candidate arises as a bound state. Although it is conceivable that the new fermions couple to the SM sector only through gravitational interactions, here we will consider the possibility that they are charged under the SM gauge group. This choice in general implies easier experimental signatures from the dark sector, but can lead at the same time to strong constraints from electroweak precision tests if the dark dynamics breaks spontaneously the SM gauge symmetry. It is thus interesting to investigate under which conditions the new dynamics preserves the SM gauge group.

2 New dynamics and new symmetries

The new dark dynamics can be broadly characterized depending on whether it generates dynamically its mass scales and whether it preserves the SM gauge symmetry.

Consider a theory defined in terms of a set of left-handed Weyl fermions Ψ transforming under the gauge group G as the direct sum of irreducible, finite-dimensional, unitary representations, $\Psi = \oplus_k \psi^{(r_k)}$. It follows:

Theorem 1 (mass term):

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[1]Stability of the DM as a consequence of an accidental symmetry was considered by many authors and appeared, for example, in the early work on technibaryon dark matter of Ref. [5)]. It was recently emphasized and its consequences were thoroughly spelled out in Ref. [6)].

[2]Additional fundamental scalars are also an option but one should make sure that they do not reintroduce hierarchy problems which cannot be addressed by the same mechanisms at work for the Higgs field. One could also envisage even more economical scenarios with only new dynamics or only new matter. In the case with only new dynamics where the DM candidate is a dark glueball, the relic density can be correctly reproduced if the dark and SM sectors are decoupled throughout their thermal histories [7, 8)]. The case with only new matter is realized by theories of Minimal Dark Matter [9)]. The new matter must come in real or pseudo-real representations of the SM gauge group and have a sufficiently large bare mass. Such DM scale is not determined dynamically by the theory, but it could be so by extending the latter to a more fundamental description at high energy where matter fields fill complex representations.

A mass term $\psi^{(r_j)}\psi^{(r_i)}$ is allowed in the Lagrangian if:

(A) $r_j \sim \bar{r}_i$

and (only for $r_i = r_j$)

(B) $\psi^{(r_j)}\psi^{(r_i)}$ *is overall symmetric in gauge and flavor space.*

Condition *(A)* requires that r_j be unitary equivalent to the conjugate representation $\bar{r}_i$, i.e. that there exists a unitary matrix S such that:

$$S^{-1}U(r_j)S = U^*(r_i)\,, \qquad SS^\dagger = 1\,. \tag{1}$$

In order for this to hold, of course, r_i and r_j must have the same dimension. It is a simple result of group theory that the product $r_i \times r_j$ contains a singlet of G, as required in order to write a mass term, if and only if *(A)* is verified. It is also easy to show that, for irreducible r_i and r_j, the matrix S is unique up to an overall phase [10].

For $r_i = r_j$, condition *(B)* ensures that a bilinear $\psi^{(r_j)}\psi^{(r_i)}$ singlet under the Lorentz group is overall antisymmetric under the exchange of the two fermion fields, as required for a mass term. In this case it follows that $SS^t = \pm 1$ and one distinguishes between two possibilities: [10]

for $S = S^t$ — r_i is called real-positive, or real. There exists a unitary transformation R which makes the generators of $U(r_i)$ purely imaginary and antisymmetric, with $RR^t = S$;

for $S = -S^t$ — r_i is called real-negative, or pseudo-real.

A second useful result which holds true is the following:

Theorem 2 (condensate):
A scalar condensate $\langle\psi^{(r_j)}\psi^{(r_i)}\rangle$ can be a singlet of the gauge symmetry group, and can thus preserve it, only if conditions (A) and (B) are satisfied.

Energy considerations suggest that if a scalar fermion condensate is allowed which is a singlet of the confining gauge group G_{strong}, it will dominate over other possible condensates and align the vacuum along a G_{strong}-preserving direction [11]. The fate of the remaining gauge symmetry, G_{weak}, in this case depends on whether $\langle\psi\psi\rangle$ can be a singlet of the whole group $G = G_{strong} \times G_{weak}$. If conditions *(A)* or *(B)* are not satisfied, then G_{weak} is necessarily broken by the condensate, as a consequence of Theorem 2. Identifing G_{strong} with the whole dark dynamics and G_{weak} with the SM gauge group, G_{SM}, implies that the dark dynamics breaks spontaneously the SM gauge symmetry. One thus obtains a Technicolor (TC) theory. In original constructions of this kind [12] there exists no elementary Higgs field and the strong dynamics is entirely responsible for the EW symmetry breaking. Variants of the TC idea have also been proposed where instead the strong (dark) dynamics induces a vev for an elementary Higgs field besides contributing to the EW scale (see for example 'Bosonic Technicolor' [13] and 'Superconformal Technicolor' [14]). It is worth mentioning that while in all the above theories it is condition *(A)* which is violated, i.e. the fermion representations are complex, it is possible to construct models where Theorem 2 fails because *(B)* does not hold. Consider for example a theory with gauge group $G = SU(N_{DC})\times SU(2)_L$, with $N_{DC} \geq 3$, and a single Weyl fermion field ψ transforming as $(adj, 4)$ of G. We assume that the dark color group $SU(N_{DC})$ confines at a scale Λ above the EW scale, and ask whether the SM electroweak group $SU(2)_L$ is spontaneously broken by the condensate $\langle\psi\psi\rangle$. In this case ψ is a pseudo-real representation r of G, since it is possible to find a unitary transformation $S = -S^t$ so that $r \sim \bar{r}$. Consequently, $\psi\psi$ can be a singlet of G, but this turns out to be antisymmetric under the exchange of the two fermions fields. Hence condition *(B)* is violated and the dark-color preserving condensate breaks $SU(2)_L$. Notice that this theory is free of global anomalies [15)].[3] In general, it is possible to show that for a pseudo-real ψ, the singlet $\psi\psi$ is always antisymmetric in the corresponding gauge indices (conversely, $\psi\psi$ is always symmetric for ψ real). Technicolor theories can thus be constructed with pseudo-real representations provided there is no global symmetry group.

The above discussion suggests that if one wants to build a model where no bare masses are allowed and where the condensate does not break the SM gauge group, there are two possible routes. The first is to consider fermion representations which are real (or pseudo-real) under G_{SM} but violate *(A)* because they are complex under G_{strong}. In this case, however, a scalar fermion condensate cannot be a singlet of G_{strong} and it is common belief that the strong dynamics is spontaneously broken, i.e. Higgsed, in the infrared (IR). Arguments based on the Most Attractive Channel criterion, for example, suggest that the theory may tumble to another one with smaller gauge group [16)], and that this process continues until one of the following situations is realized: the gauge symmetry is completely Higgsed; the theory confines; the theory flows to an Abelian phase. While the above behavior is plausible, there is currently no rigorous way to define strongly-coupled chiral theories (i.e. theories with complex representations under the strong group) in a non-perturbative way and thus determine their IR phase. Progress in our understanding of quantum field theory is thus needed before one can construct sensible phenomenological models of this kind. The second route to forbid bare masses consists in taking G_{weak} larger than the SM group and let the new fermions transform as real (or pseudo-real) representations under G_{SM} and G_{strong}, but as complex representations under G_{weak}. In this case the fermion condensate aligns along a G_{strong}-preserving direction and breaks G_{weak}. Whether it breaks G_{SM} or just the remaining part of G_{weak} is a dynamical issue, i.e. depends on the vacuum alignment. If one considers the strong dynamics in isolation, there is a degenerate surface of differently oriented vacua. The degeneracy is lifted when interactions weaker than G_{strong}, in particular those of G_{weak}, are included.

As an example, consider a theory [17)] defined in terms of left-handed Weyl fermions charged under a dark $SU(N_{DC}) \times U(1)_{DC}$, with $N_{DC} \geq 3$, and the SM electroweak $SU(2)_L$ as follows:

	$SU(N_{DC})$	$U(1)_{DC}$	$SU(2)_L$
ψ_1	$\Box$	1	$\Box$
ψ_2	$\Box$	-1	$\Box$
$\bar{\psi}_1$	$\bar{\Box}$	$-a$	$\bar{\Box}$
$\bar{\psi}_2$	$\bar{\Box}$	a	$\bar{\Box}$

The charge a is an arbitrary number between 1 and -1. This is a simple extension of the chiral model proposed in [18)] where the new fermions are charged under the SM group. Overall the fermion representations are complex, but they are vector-like with respect to $SU(N_{DC})$ and $SU(2)_L$. We will assume that the subgroup $SU(N_{DC})$ gets strong above the EW scale and confines. The pattern of dynamical symmetry breaking is then determined by the Vafa-Witten theorem [19)], similarly to the QCD case: The $SU(4)\times SU(4)$ global symmetry acting on the ψ_i and $\bar{\psi}_i^*$ is spontaneously broken to the diagonal $SU(4)_d$.

[3]If ψ tranformed as $(adj, 2)$ of $SU(N_{DC}) \times SU(2)_L$, i.e. as a spin 1/2 instead of spin 3/2 of $SU(2)_L$, the theory would have a global anomaly.

For $a > 0$, $SU(2)_L$ is contained in $SU(4)_d$ and is thus preserved, while $U(1)_{DC}$ is non-linearly realized. The phenomenology of this model, including aspects related to the DM composition and relic density, will be discussed elsewhere [17].

An important class of theories where the condensate instead aligns along a G_{SM}-breaking direction is that of Composite Higgs (CH) models [20, 21, 22]. In these constructions one requires that the set of Nambu-Goldstone bosons (NGBs) which arise from the global symmetry breaking induced by the strong dynamics includes an $SU(2)_L$ doublet. The latter plays the role of the composite Higgs field. Electroweak symmetry breaking thus follows dynamically from vacuum misalignment, but it can be also described conveniently as a two-step process: first, at some scale f the strong dynamics confines generating the NGBs; these then acquire a potential from weaker radiative corrections and trigger EWSB at a scale $v < f$. As an example consider the model by Dugan, Georgi and Kaplan [20] defined in terms of a gauge group $G_{strong} \times U(1)_A \times SU(2)_L \times U(1)_Y$ and five Weyl fermions transforming as

	G_{strong}	$U(1)_A$	$SU(2)_L$	$U(1)_Y$
ψ	r	$1/\sqrt{20}$	2	$1/2$
$\tilde{\psi}$	r	$1/\sqrt{20}$	2	$-1/2$
ψ_s	r	$-4/\sqrt{20}$	1	0

Here r is some real representation of G_{strong}, and the dark group includes an Abelian factor $U(1)_A$. The representations are overall complex though (pseudo-)real under G_{strong} and $SU(2)_L \times U(1)_Y$. The vacuum alignment can be determined by studying the effective potential of the composite Higgs doublet $H \sim (\psi\psi_s)$. For $5g_A^2 > 3g^2 + g'^2$ the condensate is forced to align in an $U(1)_A$-preserving direction and thus breaks $SU(2)_L \times U(1)_Y$. More in general, other CH theories have been constructed were vacuum misalignment is induced by fermion couplings, in particular interactions involving the top quark (see for example [23, 24]).

In order to get a complete classification of possibilities, it is worth analyzing also theories where the premises of Theorem 1 are fullfilled. In such case bare mass terms are allowed for the dark fermions, which thus introduce arbitrary new scales into the theory. Although constructions of this kind seem to go beyond the paradigm of minimality of the SM, they can still be relevant for our discussion if they are considered as effective theoretical layers valid up to some cutoff energy where they are embedded into a more fundamental description with complex representations. In this way, the mass scales introduced in the effective theory can be derived in terms of couplings of the UV description. Interestingly, the same situation is realized within the SM: below the EW scale the matter content forms vector-like representations of the unbroken $SU(3)_c \times U(1)_{em}$ gauge group, and quarks and leptons have a spectrum of bare masses. Above the EW scale, on the other hand, the theory has complex representations and the value of each mass is explained in terms of the corresponding Yukawa coupling.

If Theorem 1 implies the existence of bare mass terms, then from Theorem 2 it follows that the condensate can preserve the SM gauge symmetry. Whether this actually occurs or not is again a dynamical issue and depends on the vacuum alignment. The same considerations made above apply to this case as well, and one can for example construct CH models where misalignment occurs because of fermion interactions. Interesting alternative scenarios in this case are theories where vacuum misalignment originates from the mixing of the composite Higgs with an elementary one [20, 25, 26]. This is possible if Yukawa couplings are allowed between the dark fermions and the elementary Higgs. The lighter physical Higgs scalar in the spectrum is a mixture of the elementary and the composite fields and for this reason such scenario has been dubbed Partial Higgs Compositeness. See the talk by Michele Redi at this conference for more details on these theories.

Consider finally a situation in which no vacuum misalignment is generated, for example because no additional gauge interactions exist and Yukawa couplings are not allowed. In this case the dark dynamics does not play a direct role in EWSB and consequently such theories are less constrained by EW precision tests. Since bare mass terms for fermions are allowed, the scale at which new physics states first appear is arbitrary. If some of the new particles are within the reach of the LHC or future colliders, these theories can lead to interesting phenomenology and experimental signatures. For a recent study of theories with vector-like dark fermions, a scenario dubbed 'Vectorlike Confinement', see for example Ref. [27]. While the dark dynamics has no impact on EWSB, it may however play a role in explaining the observed DM abundance. If this is the case, then the value of the new physics scale can be fixed or significantly constrained. Most interesting for our discussion are the so-called scenarios of Accidental Composite Dark Matter, where the DM candidate is a composite state of the dark dynamics and its stability is a consequence of accidental symmetries. Ref. [6] performed a systematic classification of such theories focusing on $SU(N)_D$ and $SO(N)_D$ gauge groups with vector-like fermions in the fundamental representation. A robust and viable candidate of DM is given by dark baryons. For example, if the dark quarks are ligher than the confinement scale Λ, then the observed relic density is correctly reproduced for Λ of order 100 TeV. A smaller scale is instead obtained if the quark masses are heavier than Λ [28], because in this case the dark baryons are perturbative bound states and their annihilation cross section decreases. Experimental signatures of these scenarios in general come from direct detection experiments, through the relatively large electric dipole moment predicted for the dark baryons, and from the production of the lighter dark mesons at the LHC or at future colliders.

While dark baryons are a motivated possibility, other bound states of the strong dynamics can play the role of DM candidate. For example, dark mesons can be stable because of accidental species number or G-parity [6]. As a more exotic possibility, consider bound states made of one dark quark Q plus dark glue, in the case in which Q is a Weyl fermion in the adjoint of a dark $SU(N)_{DC}$. Dark Matter candidates of this kind were considered in the context of SUSY theories and dubbed 'glueballinos', as they are the partners of the glueballs [7]. Here we want to briefly discuss the case in which dark quarks have non-trivial quantum numbers also under the SM electroweak group, and will denote the corresponding bound state as the 'gluequark' [29]. The accidental symmetry which makes gluequarks stable is a Z_2 parity under which Q is odd. The spectrum thus divides into even states (glueballs, QQ mesons, etc.) and odd states (gluequarks, QQQ fermions, etc). Requiring to have an EM neutral gluequark in the spectrum, to play the role of DM, and restricting to real or vector-like representations under the SM selects a few viable quantum number assignments. The most minimal non-trivial model consists of three Weyl fermions in the adjoint of $SU(N)_{DC}$ transforming as a triplet of $SU(2)_L$ with zero hypercharge. The lowest-dimensional operator violating the accidental Z_2 in this case is $H\sigma^i \ell \sigma^{\mu\nu} Q^i G_{\mu\nu}$, with dimension 6. The spectrum includes a gluequark V^i triplet of $SU(2)_L$, whose neutral component V^0 is the DM candidate. The thermal history of the universe and the phenomenology at low energies have some distinctive features compared to models with baryonic DM. For example, gluequark DM is expected to have very small electric dipole moment, since its constituents are electrically neutral. This makes it elusive at direct detection experiments. Furthermore, in the limit of large quark masses, the gluequark is a heavy but sizable bound state: its mass is of order M_Q while its size is $\sim 1/\Lambda >> 1/M_Q$. This is to be constrasted with dark baryons, which are small –hence perturbative– bound states of radius $\sim 1/M_Q$. Non-perturbative annihilation processes have thus a much larger cross section in the case of the gluequark and can boost the value of M_Q which reproduces the DM relic density. A detailed analysis of the thermal history of the universe in this model will be presented in a forthcoming publication [29]. In general, models with gluequark DM tend to reproduce the observed relic density for larger bound state masses, and are more difficult to be discovered. In particular, in the $M_Q > \Lambda$ limit detection may come from indirect DM searches or even gravity wave signals produced during a first-order dark phase transition, rather than from collider signatures. For $M_Q < \Lambda$ the relic density is reproduced for $\Lambda \sim 50\,\text{TeV}$ and discovery may come first from the production of dark mesons at colliders, similarly to the case of baryonic DM with light quarks.

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