

Topcolor Models and Scalar Spectrum

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ABSTRACT

We review the motivation and main aspects of Topcolor models with emphasis on the spectrum of relatively light scalars and pseudo-scalars.

I. DYNAMICAL GENERATION OF m_t .

The generation of a large fermion mass like m_t is an extremely difficult problem in theories of dynamical Electroweak Symmetry Breaking (ESB). For instance, in Technicolor theories an Extended Technicolor (ETC) interaction is required in order to obtain fermion masses. The interaction of the ETC gauge bosons with fermions and technifermions gives rise to fermion masses through terms of the form

$$m_f = \frac{g_{ETC}^2}{M_{ETC}^2} \langle \bar{T}_L T_R \rangle \quad (1)$$

where M_{ETC} is the ETC gauge boson mass and $\langle \bar{T}_L T_R \rangle$ is the technifermion condensate. Thus, in order to generate the correct value of m_t , the ETC scale has to be of $\mathcal{O}(1 \text{ TeV})$, which is uncomfortably low. Several modifications of the dynamics within technicolor have been proposed in order to accommodate such a large fermion mass [1]. On the other hand, in top-condensation models, the large top-quark mass is obtained from a $\langle \bar{t}_L t_R \rangle$ arising as a consequence of a new gauge interaction, Topcolor, which couples strongly to the top quark. Topcolor generates four-fermion interactions of the form

$$\frac{g^2}{\Lambda^2} \bar{\psi}_L t_R \bar{t}_R \psi_L \quad (2)$$

where ψ is the (t, b) $SU(2)$ doublet and Λ is the typical scale of the new interactions. If the coupling g is strong enough the top condensation occurs. The top chiral symmetry is spontaneously broken and a large m_t is generated. This implies the presence of Goldstone bosons. Originally [2], it was proposed that these were identified with the longitudinal components of the electroweak gauge bosons so that Topcolor would also be fully responsible for ESB. With the ESB scale defined as $v \approx 246 \text{ GeV}$, the decay constant of these top-pions is given by the Pagels-Stokar formula

$$f_{\pi_t}^2 \simeq \frac{N_c}{16\pi^2} m_t^2 \ln \frac{\Lambda^2}{m_t^2} \quad (3)$$

with N_c the number of colors. From (3) it can be seen that, in order for $f_{\pi_t} = v/\sqrt{2}$ and m_t to be close to the measured value, the Topcolor scale Λ has to be extremely large ($\simeq 10^{15} \text{ GeV}$). This translates into an acute fine-tuning of the coupling g in (1), which has to be adjusted to the critical value with unnaturally

high precision. One way to avoid this problem within Topcolor models is to give up the idea that ESB is fully driven by the $\langle \bar{t}t \rangle$ condensate. For instance, a cutoff scale $\Lambda \simeq \mathcal{O}(1 \text{ TeV})$ gives a non fine-tuned coupling (a few percent above critical), but a top-pion decay constant of the order of $f_{\pi_t} \simeq (50 - 60) \text{ GeV}$, which gives only small masses to the W and the Z . In this version of Topcolor [3], a separate mechanism must be invoked to generate most of the W and Z masses. This is the case of Topcolor-assisted Technicolor. Most of m_W and m_Z , as well as *small* ($\simeq 1 \text{ GeV}$) quark masses come from the Technicolor sector. Thus a small portion of m_t also comes from ETC terms, but most of the top-quark mass is dynamically generated by the Topcolor mechanism. The explicit ETC quark mass terms for top and bottom turn the top-pions into massive pseudo-Goldstone bosons. Their masses can be estimated in the fermion loop approximation to be [3]

$$m_{\pi_t}^2 \simeq \frac{N_c}{8\pi^2} \frac{m_{ETC} m_t}{f_{\pi_t}^2} \Lambda^2 \quad (4)$$

where m_{ETC} is the effective value of the ETC quark masses. Although these are initially of the order of 1 GeV , the ETC top-quark mass receives large radiative enhancements from the Topcolor interactions [3]. Thus one can have top-pion masses in the range $m_{\pi_t} \simeq (100 - 300) \text{ GeV}$. If $m_{\pi_t} < m_t$, then the top quark would primarily decay as $t \rightarrow \pi_t^+ b$. The current CDF measurement of $Br(t \rightarrow W^+ b)$ implies $m_{\pi_t} > 150 \text{ GeV}$ at 68% confidence level [4]. In what follows we will assume that the Topcolor enhancement to the ETC top mass is enough to make $m_{\pi_t} > m_t$.

The existence of the top-pions ($\pi_t^\pm, \pi_t^0$) is an essential ingredient in the Topcolor scenario, regardless of the dynamics responsible for the ESB sector and the other quark masses. The presence of other low-lying states depends on the details of the model. A Topcolor model is greatly specified by choosing a mechanism of isospin breaking, which selects the top-quark direction for condensation, leaving the bottom quark unaffected. The complete anomaly-free fermion content is necessary in order to know the scalar spectrum of the model. These low lying states have, in most cases, masses well below the cutoff scale Λ , which points at them as possibly the first signal for Topcolor.

II. TOPCOLOR MODELS AND SCALAR SPECTRUM

In all models the Topcolor group contains an $SU(3)_1 \times SU(3)_2$ which at an energy scale Λ breaks down to ordinary $SU(3)_c$. The $SU(3)_1$ is assumed to interact strongly with the third generation quarks. After Topcolor breaking there is, in addition to the massless gluons, an octet of massive colored vector particles: the top-gluons. At this stage, if the Topcolor coupling is above

critical we would have both t and b condensation. To avoid the latter, the effective Topcolor interaction must be isospin breaking. We describe two typical scenarios to implement this aspect of the theory.

A. Models with an additional $U(1)$

An effectively isospin breaking interaction is obtained by embedding two new $U(1)$ interactions in the weak hypercharge group such that $U(1)_1 \times U(1)_2 \longrightarrow U(1)_Y$, with the $U(1)_1$ strongly coupled to the third generation. This leaves an additional color-singlet massive vector boson, a Z' . Both the top-gluon and the Z' have masses of order Λ . After integrating out these heavy particles, the interesting effective four-fermion interactions that are induced have the form [8]

$$\mathcal{L} = \frac{4\pi}{M_B^2} \left\{ \left(\kappa + \frac{2\kappa_1}{9N_c} \right) \bar{\psi}_L t_R \bar{t}_R \psi_L + \left(\kappa - \frac{\kappa_1}{9N_c} \right) \bar{\psi}_L b_R \bar{b}_R \psi_L \right\} \quad (5)$$

Here, $\kappa = (g_3^2/4\pi) \cot^2 \theta$ and $\kappa_1 = (g_1^2/4\pi) \cot^2 \theta'$, with g_3 and g_1 the QCD and $U(1)_Y$ couplings respectively. The angles θ and θ' characterize the embedding of the Topcolor and the $U(1)_1 \times U(1)_2$ groups in $SU(3)_c$ and $U(1)_Y$. The requirement that $SU(3)_1$ and $U(1)_1$ couple strongly to the third generation translates into the conditions $\cot^2 \theta \gg 1$, $\cot^2 \theta' \gg 1$. The criticality condition

$$\kappa - \frac{\kappa_1}{9N_c} < \kappa_{\text{critical}} < \kappa + \frac{2\kappa_1}{9N_c} \quad (6)$$

must be satisfied in order to obtain $\langle \bar{t}t \rangle \neq 0$ and $\langle \bar{b}b \rangle = 0$. Constraints on the top-gluon sector come from t production, as well as t and b dijet mass distributions at the Tevatron [5]. This specific model is also constrained by the effects of the Z' on low energy data, both at the Z pole [6] and at low energies through FCNC [7, 8]. However, it is possible to accommodate all these constraints with a Topcolor scale $\Lambda \gtrsim 1$ TeV and still not have a fine tuning problem. For instance, in the most general case, a 2 TeV top-gluon gives a Topcolor coupling 4% above its critical value. Therefore, one can imagine a scenario where the Topcolor gauge bosons are at the few-TeV scale, making their direct detection difficult. In this scenario it is possible that the effects of Topcolor dynamics will appear at lower energy scales due to the presence of a relatively light scalar spectrum.

The effective interactions of (5) can be written in terms of two auxiliary scalar doublets ϕ_1 and ϕ_2 . Their couplings to quarks are given by [9]

$$\mathcal{L}_{\text{eff}} = \lambda_1 \bar{\psi}_L \phi_1 t_R + \lambda_2 \bar{\psi}_L \phi_2 b_R + \text{h.c.} \quad (7)$$

where $\lambda_1^2 \equiv 4\pi(\kappa + 2\kappa_1/9N_c)$ and $\lambda_2^2 \equiv 4\pi(\kappa - \kappa_1/9N_c)$. At energies below Λ the auxiliary fields acquire kinetic terms, becoming physical degrees of freedom. With the properly renormalized fields $\phi_i^r = z_i^{1/2} \phi_i$ the criticality conditions (6) are equivalent to

$$\langle \phi_1^r \rangle = f_{\pi_t}, \quad \langle \phi_2^r \rangle = 0 \quad (8)$$

The ϕ_1 doublet acquires a Vacuum Expectation Value (VEV) giving mass to the top quark through the coupling in (7). It is of the form

$$\phi_1^r = \begin{pmatrix} f_{\pi_t} + \frac{1}{\sqrt{2}}(h_t + i\pi_t^0) \\ \pi_t^- \end{pmatrix} \quad (9)$$

As mentioned earlier, the set of three top-pions acquires a mass from explicit small quark mass terms (the ETC masses in Topcolor-assisted Technicolor). There is also a scalar, the h_t or top-Higgs, which mass is estimated in the Nambu–Jona-Lasinio (NJL) approximation to be

$$m_h \simeq 2m_t \quad (10)$$

The second doublet is present as long as the Topcolor interaction couples to b_R , as is the case in (5). It is given by

$$\phi_2^r = \begin{pmatrix} H^+ \\ \frac{1}{\sqrt{2}}(H^0 + iA^0) \end{pmatrix} \quad (11)$$

These states are deeply bound by the Topcolor interactions and therefore can be light. Their masses can be estimated once again within the NJL approximation. For instance for $\kappa_1 \simeq 1$ and $\Lambda = (2 - 3)$ TeV one has, for the neutral states [9],

$$m_{H,A} \simeq (150 - 330) \text{ GeV}, \quad (12)$$

whereas the mass of the charged states is determined by the relation

$$m_{H^\pm}^2 = m_{H,A}^2 + 2m_t^2. \quad (13)$$

The couplings to quarks can be read off equation (7). Calculating the field renormalization constants z_i to one loop and using (3), the couplings are simply m_t/f_{π_t} where m_t is the dynamically generated top quark mass. This is a typical Goldberger-Treiman factor. Recalling that in Topcolor models we expect $f_{\pi_t}/v \simeq 3$, we see that the coupling of the Topcolor “Higgs” sector to the top quark is considerably larger than that of the SM Higgs boson. Moreover, in models where the second doublet is present this couples to b quarks with the same strength as (h_t, π_t^0) couple to top. This implies that H^0 and A^0 decays are dominated by the $b\bar{b}$ final state. The existence of these relatively light scalar states strongly coupled to third generation quarks implies a very rich phenomenology. In what follows we analyze the implications of the scalar spectrum in the model described above. We discuss other alternatives in model building in the next section, with emphasis on the differences in the scalar spectrum and phenomenology.

Top-pions: As discussed earlier, these are the pseudo-Goldstone bosons of the breaking of the top chiral symmetry. They couple to the third generation quarks as

$$\frac{m_t}{\sqrt{2}f_{\pi_t}} (i\bar{t}\gamma_5 t \pi^0 + i\bar{t}_R b_L \pi^+ + i\bar{b}_L t_R \pi^-) \quad (14)$$

Although their masses are lifted by the ETC interactions, they can still play an important role in low energy observables such as as rare B decay branching fractions and angular distributions [7, 8], as well as in electroweak precision observables at the Z pole like R_b [10]. Top-pions do not have two-gauge-boson

couplings, and thus single π_t production must involve a triangle diagram. At hadron colliders they are mostly produced through the gluon-gluon- π_t^0 effective coupling induced by the top loop. For instance, the gluon-gluon fusion has a cross section larger than that of the s-channel production of the SM Higgs boson by a factor of [11]

$$r^2 \equiv \left(\frac{v}{\sqrt{2}f_{\pi_t}} \right)^2. \quad (15)$$

This enhancement is also present in the π_t production in association with top quarks, or in any production mechanism involving the $\bar{t}t\pi_t^0$ coupling. Production at e^+e^- colliders is discussed in [11].

If $m_{\pi_t} < 2m_t$ the top-pion would be narrow ($\Gamma_{\pi_t} \lesssim 1$ GeV). In models as the one presented above, where b_R couples to the strong $SU(3)_1$ interaction, instanton effects induce a coupling of π_t to b_R , which is not present in (14). This implies that, as long as the $\bar{t}t$ channel is not open, the dominant decay mode of the π_t^0 is to $b\bar{b}$ [11]. Also, and independently of the presence of instanton effects, there is a coupling of $\bar{c}c$ to π^0 given by (14) times two powers of the $t \rightarrow c$ mixing factor arising from the rotation of weak to mass quark eigenstates. However, in the present model the $b\bar{b}$ mode is expected to dominate.

Top-Higgs: It corresponds to a loosely bound, CP even, $\bar{t}t$ state, coupling to the top quark with strength $m_t/(\sqrt{2}f_\pi)$. In the NJL approximation its mass is given by (10), and therefore details of the non-perturbative dynamics are crucial to understand the decay modes and width of the h_t . The top-Higgs production is analogous to the π_t case. The main difference between these two states is that h_t couples to gauge boson pairs. These couplings are *suppressed* with respect to the case of a SM Higgs by a factor of $1/r$. However, if the $\bar{t}t$ channel is not open this would be the dominant decay mode. If this is the case, Γ_{h_t} will be considerably smaller than the width of a SM Higgs, whereas its production cross section will still be r^2 times larger. On the other hand, if $m_{h_t} > 2m_t$ the cross section is still the same but Γ_{h_t} is r^2 larger than the width of the SM Higgs of the same mass, so that h_t could only be detected as an excess in the $\bar{t}t$ in a given channel.

The “b-pions” $H_b^0, A_b^0, H_b^\pm$: If these states are present (i.e. if Topcolor couples to b_R) they give potentially dangerous contributions to various low energy observables. Their couplings to quarks are analogous to those in (14). The strongest constraint on a model containing these bound-states comes from $B^0 - \bar{B}^0$ mixing [9] and implies the existence of large suppression factors in the quark mixing matrices [8]. A constraint independent of these details is the contribution to $\Delta\rho_* = \alpha T$, the parameter measuring deviations from the SM ρ parameter, due to the splitting between the neutral and charged states. As an illustration of the size of the effect, we plot $\Delta\rho_*$ in Fig. 1 as a function of the mass of the neutral scalars and making use of the NJL result (13) for the mass splitting. The horizontal lines represent the 95% c.l. interval obtained using $\alpha_s(M_Z) = 0.115$ [12]. Although the bound is tighter as $\alpha_s(M_Z)$ increases, the b-pion splitting, always present in models with the additional $U(1)$ ’s, is not in contradiction with electroweak precision measurements.

These states decay almost exclusively to b pairs and they are

extremely broad. Their production proceeds in similar ways to that of h_t and π_t , with the important difference that the quark inside the triangle loop, the b quark, is much lighter than $\sqrt{s}$ which translates into an effective suppression of the amplitude from its “hard” value of m_t/f_{π_t} .

In addition to these low lying scalar states, the Topcolor interactions in principle lead to the formation of heavier bound states. For instance, there will be a color singlet vector meson, the top-rho ρ_t . However, ρ_t not only couples to top-pions but also couples directly to third generation quarks and with strength proportional to m_t/f_{π_t} . Moreover, their mass can be estimated in the NJL approximation to be of the order of the cutoff Λ [10]. This suggests that the influence of ρ_t in low energy observables as well as in production of Topcolor bound-states (e.g. top-pions through vector meson dominance) is largely suppressed.

B. “Axial” Topcolor

In this type of Topcolor models, the strong $SU(3)_1$ group does not couple to b_R , barring the possibility of a $\langle b_L b_R \rangle$ condensate. An example was presented in [8]. The cancelation of anomalies requires the introduction of a new set of fermion fields, $Q_{L,R}^a$, with $a = 1, \dots, N_Q$. There is a new interaction, $SU(N_Q)$. The novelty is that, depending on the choice of N_Q , light quarks might have to “feel” the strong Topcolor interaction. For instance for $N_Q = 3$ the fermions must transform under $SU(3)_Q \times SU(3)_1 \times SU(3)_2$ as

$$\begin{aligned} (t, b)_L (c, s)_L &\simeq (1, 3, 1) \\ t_R &\simeq (1, 3, 1) \\ Q_R &\simeq (3, 3, 1) \\ (u, d)_L &\simeq (1, 1, 3) \\ (u, d)_R (c, s)_R &\simeq (1, 1, 3) \\ b_R &\simeq (1, 1, 3) \\ Q_L &\simeq (3, 1, 3). \end{aligned}$$

As advertised above, b_R is not coupled to the strong Topcolor interaction, whereas the cancelation of anomalies now requires that $(c, s)_L$ transforms as a triplet under $SU(3)_1$. The

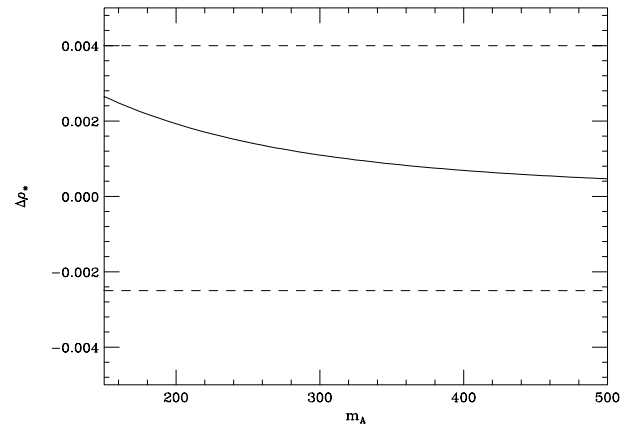


Figure 1: The contribution to $\Delta\rho_*$ due to the mass splitting among b-pions.

quarks have standard $U(1)_Y$ assignments, whereas the $Q_{L,R}$ have $Y = 0$. Furthermore, they are $SU(2)_L$ singlets and therefore electrically neutral. If leptons are incorporated with their standard $SU(2)_\times U(1)_Y$ quantum numbers and as singlets under $SU(3)_Q \times SU(3)_1 \times SU(3)_2$, all anomalies cancel. The $SU(3)_Q$ forms a $\langle Q\bar{Q} \rangle$ condensate which, in turn, breaks Topcolor down to $SU(3)_c$ dynamically.

The $SU(3)_1$ is chiral-critical and leads to the formation of a $t\bar{t}$ condensate and a dynamical m_t . This breaks an $SU(4)_L \times U(1)_L \times U(1)_R$ global chiral symmetry and leads to the existence of a composite scalar field F , quadruplet under $SU(4)_L$. This can be decomposed into two doublets, one of which is just ϕ_1^+ of Section II.A, containing h_t and the top-pions. The additional doublet, $\mathcal{C}$, contains a set of three pseudo-scalars, the “charm-top-pions” or $\vec{\pi}_c$, as well as a “charm-top-Higgs” h_c :

$$\mathcal{C} = \begin{pmatrix} \frac{1}{\sqrt{2}}(h_c + i\pi_c) \\ \pi_c^- \end{pmatrix} \quad (16)$$

The masses of the charm-top-pions are generated by the same terms that break chiral symmetry explicitly and induce the top-pion masses, so they are expected to be of the same order as m_{π_t} . The mass of the h_c in the NJL approximation is $\approx m_t$. The couplings to quarks are analogous to those of the top-pions and have the form [8]

$$\frac{m_t}{f_{\pi_t}} (\bar{c} \bar{s})_L t_R \mathcal{C} + \text{h.c.} \quad (17)$$

Important constraints on this scenario come from low energy observables. For instance, $D^0 - \bar{D}^0$ mixing is mediated by three-level s-channel exchange of h_c and π_c^0 . However, the mixing amplitude is proportional to the product of the various unknown quark rotation factors entering the transformation from the weak to the mass eigen-basis for quarks [8]. These factors depend on the sector of the theory responsible for light quark masses, e.g. ETC. Thus, charm mixing is a direct constraint on the light-quark mass matrices generated by this sector. On the other hand, charm-top-pion contributions to R_c and R_s are potentially large [10] and independent of these details.

The production of top-pions in this model is completely analogous to the model in the previous section. However, here the top-gluons do not couple to b_R , so there will be no instanton induced b quark mass term and therefore no $\pi_t^0 \rightarrow b\bar{b}$ decay mode. Thus, if top-pions have a mass below the $t\bar{t}$ threshold, the neutral states will primarily decay to gluons. Also and as we mentioned in the previous model, the $c\bar{c}$ decay mode is suppressed by two factors coming from the $t \rightarrow c$ rotations. Although these factors are not very constrained, naive estimates lead to values of $\Gamma(\pi_c^0 \rightarrow c\bar{c})$ that indicate that this mode is competitive with the gluon-gluon channel [11].

The production of charm-top-pions does not proceed via gluon-gluon fusion. Production through the anomaly is no only suppressed by loop factors and couplings but also by CKM factors. The most efficient mechanism for single production of h_c or $\pi_c^{0,\pm}$ is by emission off a top quark line, with this turning into a charm or strange quark for neutral and charged emission respectively. The multijet final state, e.g. $t\bar{t}c\bar{c}$ for the neutral

case, may be difficult to separate from the QCD background, especially given that these scalars are expected to be very broad. More detailed studies are necessary. Charm-top-pions can also be pair produced, just like the top-pions and the top-Higgs, via their model independent couplings to gauge bosons. However, as in the case of the previous section, these couplings are also suppressed by the ratio f_{π_t}/v .

It is possible to extend this model to the first family by just including the $(u, d)_L$ among the quarks strongly coupled to the Topcolor interactions, in addition to $(c, s)_L$. Such Topcolor models are partly motivated by the possibility of the existence of an excess of events at high energies in the CDF data for the inclusive jet cross section [13]. In Ref. [14] the consequences of a universally coupled top-gluon on the inclusive jet production were studied. It is possible to write down an anomaly free realization of this proposal. The only other modification needed to insure the cancelation of all anomalies is that now we need $N_Q = 5$. Thus, $(u, d)_L$ transforms as $(1, 3, 1)$ under $SU(5)_Q \times SU(3)_1 \times SU(3)_2$, and the other fermions transform as before, with the exception of the $Q_{L,R}$ that transform as quintuplets under $SU(5)_Q$. The global chiral symmetry broken by the dynamical top quark mass is now $SU(6)_L \times U(1)_R$. Therefore, besides the scalar content of the two previous models, there is a new scalar doublet, $\mathcal{U}$, the “up-top-pions”. Their masses are very similar to those of h_t and m_{π_c} . Their couplings to quarks can be read off (17), with the replacements $\mathcal{C} \rightarrow \mathcal{U}$ and $(\bar{c} \bar{s})_L \rightarrow (\bar{u} \bar{d})_L$. Although these couplings imply an additional contribution to $D^0 - \bar{D}^0$ mixing, it is governed by the same product of up-quark rotation matrix elements, and therefore can be avoided by the same choice of mass matrices that suppressed the charm-top-pion contributions. There will also be new contributions to the hadronic Z width [10]. The production of h_u and π_u has the same features as that of the charm-top-pions. In both cases, one expects these resonances to be very broad and a careful study is needed in order to establish their detectability in the various channels and in different environments.

III. CONCLUSIONS

We have reviewed the main aspects of Topcolor models and the constraints derived from the spectrum of low laying scalars and pseudo-scalars. A basic feature of all Topcolor models is the existence of the loosely bound state h_t and a triplet of top-pions $\vec{\pi}_t$. The latter is present in the physical spectrum in models where Topcolor is not solely responsible for ESB (e.g. Topcolor-assisted Technicolor).

Although the production of h_t shares several features with that of the SM Higgs, it has a larger cross section and a very different width (larger or smaller by r^2 depending on m_{h_t}). On the other hand, the observation of π_t^0 at hadron colliders is problematic given that its decay modes, $b\bar{b}$, $c\bar{c}$ and gluon-gluon, are very hard to extract from the background. Its production at lepton colliders is discussed in [11].

Specific realizations of Topcolor tend to have additional, relatively light, scalars. Given that these tend to be very broad objects, their detectability at various facilities requires a careful study, particularly at hadron colliders. Considering that Top-

color theories are still at relatively early model building stages, the study of low energy signals of and constraints on the scalar spectrum in the various models will play a central role in determining what the dominant Topcolor phenomenology will be in future experiments.

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