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The axigluon, a four-site model and the top quark forward–backward asymmetry at the Tevatron

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ABSTRACT

Very recently, it has been shown by other authors that the CDF measurement of the top quark forward–backward asymmetry can be explained by means of a heavy and broad axigluon. In order to work, this mechanism needs that the axigluon coupling to the top quark must be different than the coupling to light quarks and both must be stronger than the one predicted in classical axigluon models. In this Letter, we argue that this kind of axigluon can be accommodated in an extended chiral color model we proposed previously. Additionally, we show that the desired features can be derived from a simple four-site model with delocalized fermions.

1. Introduction

Although the High Energy Physics community is unanimous in considering the Standard Model (SM) to be incomplete for many and well-known reasons, such as the Higgs mass instability under quantum corrections or the unexplained hierarchy observed in the masses of the fermionic sector; it is impossible to deny its enormous (and, to some extent, unexpected) experimental success. Indeed, after years of scrutiny and new physics search, first at the LEP and the Tevatron and now at the LHC, the SM has not been still seriously challenged by collider experiments. The only exception today may be the 3.4σ deviation recently reported by CDF in the top-quark forward–backward asymmetry ($A_{FB}^{t\bar{t}}$) [1] (see also [2] and [3]). The fact that the process under question involves the top quark makes it even more interesting since, due to its mass which is near the electroweak scale, it has been suspected for a long time that the top-quark plays a special rôle in nature.

A natural candidate to explain a deviation in $A_{FB}^{t\bar{t}}$ is an axigluon (other alternatives beyond the SM and their consequences for the LHC have been considered in [4–16] and [17]). In fact, this is an expected axigluon signal [18]. It is true that it has been argued that the present signal cannot be explained by an axigluon due to other experimental constraints over it [19]. Nevertheless, Bai, Hewett, Kaplan and Rizzo have argued that a heavy ($m_A \geqslant 1.5$ TeV) and broad ($\Gamma/m_A \geqslant 0.2$) axigluon can be a viable explanation of the $A_{FB}^{t\bar{t}}$ while remaining invisible in the dijet spectrum [20].

In their paper, they parameterize the axigluon interaction as

$$g_s g_A \bar{\psi}\frac{\lambda^a}{2}\gamma_\mu\gamma_5 A^\mu\psi \tag{1}$$

where g_s is the usual QCD coupling constant a g_A measures the deviation of the axigluon coupling constant from QCD. In order to evade experimental constraints it proves useful to consider that g_A may take different values for light quarks (g_A^q) and for the top quark (g_A^t) [21]. The additional conditions that g_A^q and g_A^t must satisfy to explain the $A_{FB}^{t\bar{t}}$ anomaly and simultaneously evade experimental constrains can be extracted from the analysis made in [20] and can be (roughly) summarized as:

$$|g_A^q|, |g_A^t| > 1 \tag{2}$$

$$|g_A^q| < |g_A^t| \tag{3}$$

$$g_A^q g_A^t < 0 \tag{4}$$

Condition (2) is needed to have a broad axigluon and for making the asymmetry large enough. Condition (3) allows us to evade constrains coming from resonance and contact interaction searches. Finally, condition (4) is needed to obtain the correct sign for the asymmetry.

The construction of a model that can incorporate all these conditions, specially the different coupling constants for different

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generations, has proven to be difficult. As a solution, the authors of [20] propose a two-site model but they have to include exotic vector-like quarks in order to satisfy the $|g_A^q| < |g_A^t|$ constrain.

In this Letter we want to point out that the construction of a viable axigluon model (viable in the sense that it satisfies the conditions exposed above) can be achieved, without introducing new fermions, by enlarging the gauge sector of the model. In Section 2 we show that an effective model, based on $SU(3)_L \times SU(3)_R$ but with four octet fields transforming as gauge bosons, has enough structure to accommodate all the desire features of the axigluon. In Section 3, we show that the essential characteristics of the effective model can be obtained from a four-site model with delocalized fermions. Finally, in Section 4, we summarize our conclusions.

2. An effective model

We start by recalling an effective extended chiral color model we have previously proposed [22]. In this model, we consider an usual $SU(3)_L \times SU(3)_R$ with two pairs of vectors fields: l_μ and L_μ transforming as gauge fields of $SU(3)_L$, and r_μ and R_μ transforming as gauge fields of $SU(3)_R$. The gauge sector of the model (including a non-linear sigma model sector which is introduced to produce the breaking of $SU(3)_L \times SU(3)_R$ to $SU(3)_c$) is described by the following Lagrangian:

$$\begin{aligned}\mathcal{L} = &-\frac{1}{2}\operatorname{tr}\left\{G_{L\mu\nu}G_L^{\mu\nu}\right\} - \frac{1}{2}\operatorname{tr}\left\{G_{R\mu\nu}G_R^{\mu\nu}\right\} \\ &- \frac{1}{2}\operatorname{tr}\left\{\rho_{L\mu\nu}\rho_L^{\mu\nu}\right\} - \frac{1}{2}\operatorname{tr}\left\{\rho_{R\mu\nu}\rho_R^{\mu\nu}\right\} \\ &+ \frac{M^2}{g'^2}\operatorname{tr}\left\{\left(gl_\mu - g'L_\mu\right)^2\right\} + \frac{M^2}{g'^2}\operatorname{tr}\left\{\left(gr_\mu - g'R_\mu\right)^2\right\} \\ &+ \frac{f^2}{2}\operatorname{tr}\left\{D_\mu U^\dagger D^\mu U\right\} \end{aligned} \tag{5}$$

where

$$\begin{aligned} G_{L\mu\nu} &= \partial_\mu l_\nu - \partial_\nu l_\mu - ig_L[l_\mu, l_\nu] \\ G_{R\mu\nu} &= \partial_\mu r_\nu - \partial_\nu r_\mu - ig_R[r_\mu, r_\nu] \\ \rho_{L\mu\nu} &= \partial_\mu L_\nu - \partial_\nu L_\mu - ig'_L[L_\mu, L_\nu] \\ \rho_{R\mu\nu} &= \partial_\mu R_\nu - \partial_\nu R_\mu - ig'_R[R_\mu, R_\nu] \\ D_\mu U &= \partial_\mu U - ig_L l_\mu U + ig_R U r_\mu \end{aligned} \tag{6}$$

and U is a (matrix) non-linear scalar field which is introduces in order to break the original symmetry to $SU(3)_c$ through its non-zero vacuum expectation value.

In the unitary gauge ($U = \mathbb{1}$) a non-diagonal mass matrix is explicitly generated. Although the resulting mass matrix can be exactly diagonalized, in order to obtain easily manipulable expressions we will assume that $g' \gg g$ and we will write our results in the first order in g/g'. The eigenvalues, that is, the masses of the physical states are:

$$\begin{aligned} m_G &= 0 \\ m_A &= \frac{gf}{\sqrt{2}} \\ m_{G'} &= M \\ m_{A'} &= M \end{aligned} \tag{7}$$

Thus, the physical spectrum is composed of a (exactly) massless gluon, an axigluon and two degenerate (at this level of approximation) heavy gluon and axigluon. The normalized mass eigenvectors can be written, at the leading order in g/g', as [22]:

$$\begin{aligned} G_\mu &= \frac{1}{\sqrt{2}}l_\mu + \frac{1}{\sqrt{2}}r_\mu + \frac{g}{\sqrt{2}g'}L_\mu + \frac{g}{\sqrt{2}g'}R_\mu \\ A_\mu &= -\frac{1}{\sqrt{2}}l_\mu + \frac{1}{\sqrt{2}}r_\mu - \frac{g}{\sqrt{2}g'}\left(1 - \frac{m_A^2}{M^2}\right)L_\mu \\ &\quad + \frac{g}{\sqrt{2}g'}\left(1 - \frac{m_A^2}{M^2}\right)R_\mu \\ G'_\mu &= \frac{g}{\sqrt{2}g'}l_\mu + \frac{g}{\sqrt{2}g'}r_\mu - \frac{1}{\sqrt{2}}L_\mu - \frac{1}{\sqrt{2}}R_\mu \\ A'_\mu &= -\frac{g}{\sqrt{2}g'}\left(1 - \frac{m_A^2}{M^2}\right)^{-1}l_\mu + \frac{g}{\sqrt{2}g'}\left(1 - \frac{m_A^2}{M^2}\right)^{-1}r_\mu \\ &\quad + \frac{1}{\sqrt{2}}L_\mu - \frac{1}{\sqrt{2}}R_\mu \end{aligned} \tag{8}$$

Because we have now two fields that transform as gauge fields for each group (l_μ and L_μ for $SU(3)_L$ and r_μ and R_μ for $SU(3)_R$) any combination of the form $g(1-k)l_\mu + g'kL_\mu$ and $g(1-k)r_\mu + g'kR_\mu$, where k is a new arbitrary parameter, transforms as a connection of the respective algebra and hence may be used to define a covariant derivative. In principle, k can take different values in the right and left sectors but, for simplicity we will assume it is the same in both cases. With this structure it is possible to construct generalized covariant derivatives. This means that the Lagrangian describing the gauge interaction of quarks can be written as:

$$\begin{aligned}\mathcal{L} = &\frac{1}{2}g(1-k)\bar{\psi}l_\mu\gamma^\mu(1-\gamma_5)\psi + \frac{1}{2}g'k\bar{\psi}L_\mu\gamma^\mu(1-\gamma_5)\psi \\ &+ \frac{1}{2}g(1-k)\bar{\psi}r_\mu\gamma^\mu(1+\gamma_5)\psi + \frac{1}{2}g'k\bar{\psi}R_\mu\gamma^\mu(1+\gamma_5)\psi \end{aligned} \tag{9}$$

Using Lagrangian (9) and the definition of the physical fields, we can obtain the terms of the Lagrangian that couple the gluon and the axigluon to quarks

$$\mathcal{L} = g_s\bar{\psi}G_\mu\gamma^\mu\psi + g_s g_A\bar{\psi}A_\mu\gamma^\mu\gamma_5\psi \tag{10}$$

where

$$g_s \equiv \frac{g}{\sqrt{2}} \tag{11}$$

and

$$g_A \equiv 1 - \frac{m_A^2}{M^2}k \tag{12}$$

Interestingly, the mechanism described just above is related with fermion delocalization in deconstruction theory [25,26].

Now we have an axigluon with a modified coupling to quarks. Nothing in the model prevents that the value of g_A varies from one generation to another by choosing different values of k for different generations. This implies that in this framework we can easily satisfy the necessary conditions (2)–(4) (in particular, the $|g_A^q| < |g_A^t|$ constrain) for making the axigluon compatible with existing data and, at the same time, explain the $A_{FB}^{t\bar{t}}$ data from CDF.

Although this simple effective model may be enough for accommodating an axigluon with the desired properties to explain the $A_{FB}^{t\bar{t}}$ anomaly, from the theoretical point of view it would be more satisfactory to establish the model on more elegant basis. Such a structure may be provided by the so-called “Linear Moose” model in the context of deconstruction theory [23,24]. Of course, by adopting this theoretical framework, we do not intend to reproduce exactly our toy model: it is enough to reproduce the key features that make it successful in accommodating a viable axigluon. These key characteristics are: the presence of two pairs of

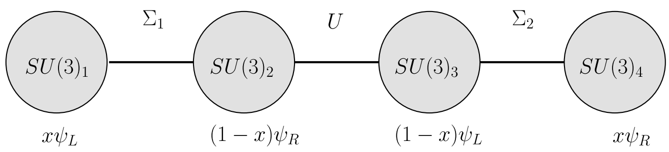

Fig. 1. Group structure of our simple four-site model.

vector bosons which transform like gauge fields and fermions delocalized in such a way that the axigluon coupling depends on the delocalization parameter. These properties can be naturally implemented in a four-site model. The construction of such a model is the subject of next section.

3. A simple four-site model

Our simple four-site model is based on the gauge group $SU(3)_1 \times SU(3)_2 \times SU(3)_3 \times SU(3)_4$ and three link fields (Σ_1, U and Σ_2) as shown in Fig. 1. We assume that the scalar fields develop vacuum expectation values ($u = \langle \Sigma_1 \rangle = \langle \Sigma_2 \rangle$ and $v = \langle U \rangle$) breaking down the gauge symmetry to $SU(3)_c$ which we identify with the usual color group.

The Lagrangian for the bosonic sector can be written as:

$$\mathcal{L} = -\frac{1}{2}\sum_{n=1}^{4} \mathrm{tr}\{F_{n\mu\nu}F_n^{\mu\nu}\} + \frac{u^2}{2}\mathrm{tr}\{D_\mu \Sigma_1^\dagger D^\mu \Sigma_1\} + \frac{v^2}{2}\mathrm{tr}\{D_\mu U^\dagger D^\mu U\} + \frac{u^2}{2}\mathrm{tr}\{D_\mu \Sigma_2^\dagger D^\mu \Sigma_2\} \tag{13}$$

where

$$\begin{aligned} F_{n\mu\nu} &= \partial_\mu A_{n\nu} - \partial_\nu A_{n\mu} - ig[A_{n\mu}, A_{n\nu}] \\ D^\mu \Sigma_1 &= \partial_\mu \Sigma_1 - igA_{1\mu}\Sigma_1 + ig\Sigma_1 A_{2\mu} \\ D^\mu U &= \partial_\mu U - igA_{2\mu}U + igUA_{3\mu} \\ D^\mu \Sigma_2 &= \partial_\mu \Sigma_2 - igA_{3\mu}\Sigma_2 + ig\Sigma_2 A_{4\mu} \end{aligned} \tag{14}$$

Notice that in order to simplify our equations we have assume that all groups share the same coupling constant g.

As usual, in the unitary gauge ($\Sigma_1 = U = \Sigma_2 = \mathbb{1}$) a non-diagonal mass matrix for the gauges bosons explicitly appears. This mass matrix can be exactly diagonalized, nevertheless, in benefice of simplicity, we work in the limit where $v \ll u$. In this limit, the physical spectrum is composed of a (exactly) massless gluon G, a light (in the sense that its mass is proportional to v) axigluon A and two degenerate heavy states (V_1 and V_2) with masses proportional to u. The original gauge fields can be written in terms of the physical fields as follows:

$$\begin{aligned} A_{1\mu} &= \frac{1}{2}G_\mu + \frac{1}{2}A_\mu + \frac{1}{2}V_{1\mu} + \frac{1}{2}V_{2\mu} \\ A_{2\mu} &= \frac{1}{2}G_\mu + \frac{1}{2}A_\mu - \frac{1}{2}V_{1\mu} - \frac{1}{2}V_{2\mu} \\ A_{3\mu} &= \frac{1}{2}G_\mu - \frac{1}{2}A_\mu - \frac{1}{2}V_{1\mu} + \frac{1}{2}V_{2\mu} \\ A_{4\mu} &= \frac{1}{2}G_\mu - \frac{1}{2}A_\mu + \frac{1}{2}V_{1\mu} - \frac{1}{2}V_{2\mu} \end{aligned} \tag{15}$$

Now we turn our attention to quarks. A crucial feature of our effective model described in the previous section was the fact that the fermions coupled to a certain linear combination of gauge bosons. In the language of deconstruction theory, it correspond to "fermion delocalization". The main idea of fermion delocalization [25,26] is that the low energy properties of the fermions (for example the color interaction of quarks) originate from the contribution of many sites of the underlying gauge group. In other words, the current couples to a linear combination of gauge fields $\sum_n x_n A_{n\mu}$ with the constrain $\sum_n x_n = 1$ in order to preserve gauge invariance under the whole group. As shown in Fig. 1, we decide to couple left-handed quarks to $A_{1\mu}$ and $A_{3\mu}$ (with weights x and $1-x$ respectively) and right-handed quarks to $A_{2\mu}$ and $A_{4\mu}$ (with weights $1-x$ and x respectively). In this way, the interaction Lagrangian for quarks can be written as:

$$\mathcal{L}_{int} = gJ^a_{L\mu}\left[xA_1^{a\mu} + (1-x)A_3^{a\mu}\right] + gJ^a_{R\mu}\left[xA_4^{a\mu} + (1-x)A_2^{a\mu}\right] \tag{16}$$

where $J^a_{L\mu}$ and $J^a_{R\mu}$ are the quark left-handed and right-handed currents, respectively. Notice that we have chosen, for simplicity, the same delocalization parameter x for right-handed and left-handed quarks. In terms of the physical fields the interaction Lagrangian can be written as:

$$\mathcal{L}_{int} = g_s\bar{\psi}\frac{\lambda^a}{2}\gamma_\mu\psi G^{a\mu} + g_s(1-2x)\bar{\psi}\frac{\lambda^a}{2}\gamma_\mu\gamma_5\psi A^{a\mu} + g_s(2x-1)\bar{\psi}\frac{\lambda^a}{2}\gamma_\mu\psi V_1^{a\mu} - g_s\bar{\psi}\frac{\lambda^a}{2}\gamma_\mu\gamma_5\psi V_2^{a\mu} \tag{17}$$

where $g_s = g/2$ is the usual strong interaction coupling constant. Our delocalization pattern has provoked the appearance of a axigluon interaction which depends on the delocalization parameter x. Indeed, in this scheme, g_A is defined as $g_A = 1 - 2x$. As in our effective model, nothing prevents that different quark generations can be differently delocalized. In this way, we can select appropriated values of x for the top quark and for the light quarks in order to consistently explain the $A_{FB}^{t\bar{t}}$ measurements. In fact, the three conditions for a viable axigluon (2)–(4) can be translated now as conditions over the delocalization parameter of the top quark (x_t) and the delocalization of light quarks (x_q). If we choose g_A^t to be negative, the x_t and x_q must satisfy:

$$\begin{aligned} &x_t > 1 \\ &x_q < 0 \\ &x_t + x_q > 1 \end{aligned} \tag{18}$$

4. Summary

In summary, we have shown that an extended chiral color model with extra color octet spin-1 fields has enough structure to describe an axigluon with appropriate coupling to quarks in such a way that it can be consistent with dijet and contact interactions searches and, at the same time, explain the $A_{FB}^{t\bar{t}}$ anomaly reported by CDF. The crucial ingredients which makes our model different from other attempts to construct a viable axigluon model, is the presence of extra gauge-like fields that allows us to implement a non-trivial fermion delocalization.

We have also shown that this scheme can be naturally realized in a four-site model. This fact is of theoretical interest since, as it is well known, deconstruction theory can be viewed as a five-dimensional theory with a discretized fifth dimension. We expect that the search for a five-dimensional origin of our model may enable us to make contact with other recently proposed explanations of the $A_{FB}^{t\bar{t}}$ anomaly [8,27].

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References

[1] T. Aaltonen, et al., CDF Collaboration, Phys. Rev. D 83 (2011) 112003, arXiv:1101.0034 [hep-ex].
[2] T. Aaltonen, et al., CDF Collaboration, Phys. Rev. Lett. 101 (2008) 202001, arXiv:0806.2472 [hep-ex].
[3] V.M. Abazov, et al., D0 Collaboration, Phys. Rev. Lett. 100 (2008) 142002, arXiv:0712.0851 [hep-ex].
[4] K.M. Patel, P. Sharma, Forward–backward asymmetry in top quark production from light colored scalars in SO(10) model, arXiv:1102.4736 [hep-ph].
[5] B. Grinstein, A.L. Kagan, M. Trott, J. Zupan, Forward–backward asymmetry in t anti-t production from flavour symmetries, arXiv:1102.3374 [hep-ph].
[6] V. Barger, W.-Y. Keung, C.-T. Yu, Tevatron asymmetry of tops in a W', Z' model, arXiv:1102.0279 [hep-ph].
[7] M.I. Gresham, I.-W. Kim, K.M. Zurek, Searching for top flavor violating resonances, arXiv:1102.0018 [hep-ph].
[8] C. Delaunay, O. Gedalia, S.J. Lee, G. Perez, E. Ponton, Extraordinary phenomenology from warped flavor triviality, arXiv:1101.2902 [hep-ph].
[9] K. Blum, C. Delaunay, O. Gedalia, Y. Hochberg, S.J. Lee, Y. Nir, G. Perez, Y. Soreq, Implications of the CDF $t\bar{t}$ forward–backward asymmetry for boosted top physics, arXiv:1102.3133 [hep-ph].
[10] B. Bhattacherjee, S.S. Biswal, D. Ghosh, Phys. Rev. D 83 (2011) 091501, arXiv:1102.0545 [hep-ph].
[11] Q.-H. Cao, D. McKeen, J.L. Rosner, G. Shaughnessy, C.E.M. Wagner, Phys. Rev. D 81 (2010) 114004, arXiv:1003.3461 [hep-ph].
[12] E.L. Berger, Q.-H. Cao, C.-R. Chen, C.S. Li, H. Zhang, Phys. Rev. Lett. 106 (2011) 201801, arXiv:1101.5625 [hep-ph].
[13] I. Dorsner, S. Fajfer, J.F. Kamenik, N. Kosnik, Phys. Rev. D 81 (2010) 055009, arXiv:0912.0972 [hep-ph].
[14] C. Delaunay, O. Gedalia, Y. Hochberg, G. Perez, Y. Soreq, Implications of the CDF $t\bar{t}$ forward–backward asymmetry for hard top physics, arXiv:1103.2297 [hep-ph].
[15] J.L. Hewett, J. Shelton, M. Spannowsky, T.M.P. Tait, M. Takeuchi, A^t_{FB} meets LHC, arXiv:1103.4618 [hep-ph].
[16] C. Degrande, J.-M. Gerard, C. Grojean, F. Maltoni, G. Servant, An effective approach to same sign top pair production at the LHC and the forward–backward asymmetry at the Tevatron, arXiv:1104.1798 [hep-ph].
[17] J.A. Aguilar-Saavedra, M. Perez-Victoria, JHEP 1105 (2011) 034, arXiv:1103.2765 [hep-ph].
[18] O. Antunano, J.H. Kuhn, G. Rodrigo, Phys. Rev. D 77 (2008) 014003, arXiv:0709.1652 [hep-ph].
[19] R.S. Chivukula, E.H. Simmons, C.-P. Yuan, Phys. Rev. D 82 (2010) 094009, arXiv:1007.0260 [hep-ph].
[20] Y. Bai, J.L. Hewett, J. Kaplan, T.G. Rizzo, JHEP 1103 (2011) 003, arXiv:1101.5203 [hep-ph].
[21] P. Ferrario, G. Rodrigo, Phys. Rev. D 80 (2009) 051701, arXiv:0906.5541 [hep-ph].
[22] A.R. Zerwekh, Eur. Phys. J. C 65 (2010) 543, arXiv:0908.3116 [hep-ph].
[23] N. Arkani-Hamed, A.G. Cohen, H. Georgi, Phys. Rev. Lett. 86 (2001) 4757, hep-th/0104005.
[24] W. Skiba, D. Tucker-Smith, Phys. Rev. D 65 (2002) 095002, hep-ph/0201056.
[25] R.S. Chivukula, E.H. Simmons, H.-J. He, M. Kurachi, M. Tanabashi, Phys. Rev. D 71 (2005) 115001, hep-ph/0502162.
[26] R.S. Chivukula, E.H. Simmons, H.-J. He, M. Kurachi, M. Tanabashi, Phys. Rev. D 72 (2005) 015008, hep-ph/0504114.
[27] C. Delaunay, O. Gedalia, S.J. Lee, G. Perez, E. Ponton, Ultra visible warped model from flavor triviality and improved naturalness, arXiv:1007.0243 [hep-ph].