The Production of Single $t$–Quarks at LEP and HERA

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Abstract

We study the possibility to produce single $t$–quarks both at LEP II and HERA. While within the Standard Model such reactions are not observable, the possibility exists in a wide class of dynamical models for the fermion mass generation. General arguments, based on hierarchical and democratic symmetries are used to arrive at $t$–production rates which are detectable.

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It is well-known that within the theoretical framework of the electroweak theory based on the gauge group $SU(2) \times U(1)$ a deeper understanding of the observed mass spectra of the leptons and quarks is not possible. The observations show that the masses in the various charged flavor channels (quarks, charged leptons) are strongly dominated by the masses of the third fermion family $m_t, m_b$ and $m_\tau$ respectively, a clear hierarchical pattern is observed. It is not known whether such a mass hierarchy is present also in the neutrino sector. In fact, it may well be that in the neutrino channel no mass hierarchy exists, which might be the reason for the large mixing angles indicated by the neutrino oscillation experiments [1].

The observed mass hierarchies and a related hierarchy of the flavor mixing angles in the quark sector suggest that nature seems to be close to the “rank–one limit”, in which all flavor mixing angles vanish and the mass matrices are proportional to the rank–one matrix [2, 3]

$$M_0 = \text{const.} \begin{pmatrix} 0 & 0 & 0 \\ 0 & 0 & 0 \\ 0 & 0 & 1 \end{pmatrix}.$$  (1)

This mass matrix can also be represented, after a linear transformation of the fermion fields, by the “democratic” mass matrix:

$$\tilde{M}_0 = \text{const.} \begin{pmatrix} 1 & 1 & 1 \\ 1 & 1 & 1 \\ 1 & 1 & 1 \end{pmatrix},$$  (2)

which exhibits an $S(3)_L \times S(3)_R$ symmetry [4, 5].

The mass generation for the members of the second and first family of fermions can be related to a specific breaking of the “democratic symmetry” [5, 6, 7].

Thus far the details of the dynamics which generate the masses and flavor mixing effects are unknown. They might be related, for example, to the presence of new interactions, or to a substructure of the leptons and quarks. We shall not speculate about those here, but rather concentrate on possible general consequences for the physics at LEP II and HERA. Once the masses of the fermions are generated, terms of the type $\bar{\psi}_L \psi_R + h.c$ ($\psi$: fermion field, $L$: lefthanded, $R$: righthanded) are introduced. These terms serve as bridges between the lefthanded fields and the righthanded ones. In the Standard Model the mass terms are, of course, provided by the interaction of the scalar field with the fermions. However, in a dynamical theory of the mass generation one expects that besides the fermion mass terms also effective interactions between the fermions and the electroweak or QCD gauge bosons are generated, which have a similar chiral structure as the mass terms, e. g. const. $\bar{\psi}_L \sigma_{\mu \nu} \psi_R F^{\mu \nu}$, where $F^{\mu \nu}$ stays generically for a field strength of an electroweak gauge boson or a gluon. Of course, such terms, if present, would have to be interpreted as signals of new interactions beyond those present in the Standard Model.
If we consider for the quarks the “rank–one” limit discussed above we can only have terms of the type \( \bar{t}_L \sigma_{\mu \nu} t_R F^{\mu \nu} \) or \( \bar{b}_L \sigma_{\mu \nu} b_R F^{\mu \nu} \) (+ h.c.). Diagonal terms like \( \bar{c}_L \sigma_{\mu \nu} c_R F^{\mu \nu} \) or flavor–changing terms like \( \bar{c}_L \sigma_{\mu \nu} t_R F^{\mu \nu} \) would violate the chiral symmetry \( SU(2)_L \times SU(2)_R \), acting on the first two families, which is present in this limit. For example, the flavor–diagonal terms for the anomalous magnetic moments can be written as

\[
\frac{e}{2\Lambda} \cdot \left( \frac{m_{0i}}{3\Lambda} \right) \bar{q}_L \sigma_{\mu \nu} N q_R F^{\mu \nu} + h.c.,
\]  

\((i = U, D, \quad m_{U0} \cong m_t, \quad m_{D0} \cong m_b)\)

\[
N = \begin{pmatrix}
1 & 1 & 1 \\
1 & 1 & 1 \\
1 & 1 & 1 
\end{pmatrix}
\]  

(4)

Here \( \Lambda \) is a scale parameter describing the energy scale of the underlying dynamics, which is supposedly of the order of \( m_t \). After diagonalization it is easily seen that the anomalous magnetic moment contributions would only be present (in the symmetry limit) for the fermions of the third family. Flavor–violating terms which could give rise to radiative decays like \( t \to c\gamma \) or \( b \to s\gamma \) do not appear.

Once this symmetry is broken and the masses of the quarks of the second family \( m_c \) and \( m_s \) enter, effective transition terms between the fields of the second and third family will in general be generated. After diagonalization of the mass terms (i. a. after rotating away flavor changing mass terms) a flavor mixing between the third and second family will appear. However, there is no reason why at the same time the flavor changing transition terms mentioned above would be rotated away. Thus we expect terms of the type \( \bar{c}_L \sigma_{\mu \nu} t_R F^{\mu \nu} \) or \( \bar{s}_L \sigma_{\mu \nu} b_R F^{\mu \nu} \) to be present. Also in the Standard Model such effective interactions are present, generated in higher orders of the electroweak interactions, but in case of the \( tc \)–transition they are too small to be interesting for experiment [8].

As an illustrative example we consider the simplest type of \( S(3)_L \times S(3)_R \)–symmetry breaking discussed in Ref. [6]. Both the mass terms of the charge \( 2/3(U) \) and charge \( -1/3 \)–quarks \((D)\) have the form:

\[
M_i = M_{0i} \begin{pmatrix}
1 & 1 & 1 \\
1 & 1 & 1 \\
1 & 1 & 1 + \varepsilon_i
\end{pmatrix}
\]  

(5)

(the index \( i \) stands for \( U, D \) respectively).

The symmetry breaking parameter \( \varepsilon \) destroys the rank one nature of the mass matrix. Thus the mass of the second family member \((c, s \) respectively) is introduced. At the same time a flavor mixing angle appears, which is proportional to the mass ratios \( m_s/m_b \) and \( m_c/m_t \) [6]. We also note that according to the observed mass spectrum \( \varepsilon_U \) is smaller than \( \varepsilon_D \): \( \varepsilon_U/\varepsilon_D \cong 0.1 \).
The anomalous magnetic moments of the quarks are described by a term proportional to

\[
\frac{e}{2\Lambda} \cdot \left( \frac{m_i}{\Lambda} \right) \bar{q}_L \sigma_{\mu\nu} \tilde{N} \ q_R \cdot F^{\mu\nu} + h.c.
\]  

(6)

where the flavor matrix \( \tilde{N} \) is given in analogy to the mass matrix by:

\[
\tilde{N}_i = \begin{pmatrix} 1 & 1 & 1 \\ 1 & 1 & 1 \\ 1 & 1 & 1 + \delta_i \end{pmatrix}
\]  

(7)

\( (i = U, D \text{ respectively}) \).

For symmetry reasons both the mass terms and the magnetic moment terms are given by matrices in family space where the \((3, 3)\)-terms are different from one. If the parameters \( \varepsilon_i \) and \( \delta_i \) were identical, the diagonalization of the mass matrices and of the magnetic moment matrices could be carried out using the same unitary transformation. Thus no flavor–nondiagonal terms would arise.

Since mass terms and magnetic moments probe different properties of the fermions, albeit they are similar in their chiral structure, there is no reason that the symmetry breaking parameters \( \varepsilon_i \) and \( \delta_i \) are the same, although they are expected to be of the same order of magnitude. Thus flavor–nondiagonal terms are expected to arise as soon as the \( S(3)_L \times S(3)_R \)–symmetry is violated. In our example these terms are proportional to \( m_c/m_t \) or \( m_s/m_b \) respectively:

\[
\mathcal{L}(t \to c) = \text{const.} \cdot \frac{e}{2\Lambda} \cdot \left( \frac{m_t}{\Lambda} \right) \left( \frac{m_c}{m_t} \right) \bar{t}_L \sigma_{\mu\nu} c_R F^{\mu\nu} + h.c.,
\]

\[
\mathcal{L}(b \to s) = \text{const.} \cdot \frac{e}{2\Lambda} \cdot \left( \frac{m_b}{\Lambda} \right) \left( \frac{m_s}{m_b} \right) \bar{b}_L \sigma_{\mu\nu} s_R F^{\mu\nu} + h.c.. \quad (8)
\]

Here the const. in front depends on how far the ratios of the symmetry breaking parameters \( \delta_i/\varepsilon_i \) differ from unity.

The example discussed above shows that the \( t \)– and \( b \)–quarks might have properties like anomalous magnetic moments, which in principle are expected to give rise to flavor–changing decays. How large these amplitudes are, depends on the expected misalignment between the mass matrices and the matrices describing the magnetic moment transitions.

From our discussion it is evident that these misalignments are due to the \( S(3)_L \times S(3)_R \)–breaking. They will not be present in the symmetry limit. In our example they are of the order of the mixing angles, i. e. of order \( m_c/m_t \) or \( m_s/m_b \). In special situations the angles describing the misalignment can be of order \( (m_c/m_t)^{1/2} \) or \( (m_c/m_b)^{1/2} \); this is the case, if the \((2,2)\)–matrix element vanishes in the magnetic moment matrix, if written in the hierarchy basis. In this case one finds e. g. for the \( t \to c \)–transition:

\[
\mathcal{L}(t \to c) = \text{const.} \cdot \frac{e}{2\Lambda} \cdot \left( \frac{m_c}{m_t} \right)^{1/2} \bar{t}_L \sigma_{\mu\nu} c_R F^{\mu\nu} + h.c. \quad (9)
\]
In order to be general we shall allow for this possibility in our discussion below.

Both for the interaction of the $t$–quark with the photon and the $Z$–boson we obtain anomalous vertices as follows:

\[
\Delta^t_Z = \bar{c} \left( i \left( C_Z + D_Z \gamma_5 \right) \sigma^{\mu\nu} \frac{q_\nu}{m_t} \right) t Z_\mu, \\
\Delta^t_\gamma = \bar{c} \left( i \left( C_\gamma + D_\gamma \gamma_5 \right) \sigma^{\mu\nu} \frac{q_\nu}{m_t} \right) t A_\mu.
\]

(10)

According to our discussion above we expect for the order of magnitudes of the vertex parameters $C, D$:

\[
\text{const.} \cdot \frac{m_c}{m_t} \cdot e < C_\gamma, \, D_\gamma < \text{const.} \cdot \sqrt{\frac{m_c}{m_t}} \cdot e \\
\text{const.} \cdot \frac{m_c}{m_t} \cdot g_Z < C_Z, \, D_Z < \text{const.} \cdot \sqrt{\frac{m_c}{m_t} \cdot g_Z}
\]

(11)

($g_Z$: $Z$–boson coupling constant)

\[
g_Z = \frac{e}{\cos \theta_W}.
\]

(12)

Thus we expect the $C, D$–parameters to be of the order of $10^{-2}$ down to the order of $10^{-3}$. Below we shall discuss the consequences for $t$–decays and $t$–production in several experimental situations.

**Top Decays**

Besides the decay $t \to b + W$ described within the Standard Model the decays $t \to c Z, \, t \to c \gamma$ proceed via the anomalous vertices. We find:

\[
\Gamma(t \to cZ) = \left( |C_Z|^2 + |D_Z|^2 \right) \frac{m_t}{8\pi} \left( 1 - \frac{M_Z^2}{m_t^2} \right)^2 \left( 1 - \frac{1}{2} \frac{M_Z^2}{m_t^2} - \frac{1}{2} \frac{M_Z^4}{m_t^4} \right),
\]

(13)

\[
\Gamma(t \to c\gamma) = \left( |C_\gamma|^2 + |D_\gamma|^2 \right) \frac{m_t}{8\pi}.
\]

(14)

The standard decay is given by:

\[
\Gamma(t \to bW) = \frac{G_F}{8\pi\sqrt{2}} |V_{tb}|^2 m_t^3 \left( 1 - \frac{M_W^2}{m_t^2} \right)^2 \left( 1 + \frac{M_W^2}{m_t^2} - 2 \frac{M_W^4}{m_t^4} \right).
\]

(15)

For the branching ratios we obtain:

\[
\text{BR}_Z \equiv \frac{\Gamma(t \to cZ)}{\Gamma(t \to bW)} = \sqrt{2} \left( 1 - \frac{M_Z^2}{m_t^2} \right)^2 \left( 1 - \frac{1}{2} \frac{M_Z^2}{m_t^2} - \frac{1}{2} \frac{M_Z^4}{m_t^4} \right) \left( |C_Z|^2 + |D_Z|^2 \right),
\]

(16)
\[ BR_\gamma \equiv \frac{\Gamma(t \to c\gamma)}{\Gamma(t \to bW)} = \frac{\sqrt{2}}{G_F |V_{tb}|^2 m_t^2} \left( \frac{1}{1 - \frac{M_W^2}{m_t^2}} \right) \left( 1 + \frac{M_W^2}{m_t^2} - 2 \frac{M_W^4}{m_t^4} \right) \left( |C_\gamma|^2 + |D_\gamma|^2 \right). \quad (17) \]

For the decays \( t \to c\gamma, t \to cZ \) exist the following experimental bounds \[ 9 \]:

\[
\begin{align*}
BR(t \to q Z) &< 0.4, \\
BR(t \to q\gamma) &< 0.029. 
\end{align*} \quad (18)
\]

Thus we obtain:

\[
\begin{align*}
|C_Z|^2 + |D_Z|^2 &< 0.16, \\
|C_\gamma|^2 + |D_\gamma|^2 &< 6.5 \cdot 10^{-3}. 
\end{align*} \quad (19)
\]

For an order of magnitude estimate we take \( C \approx D \) and find:

\[
\begin{align*}
C_Z, D_Z &< 0.3, \\
C_\gamma, D_\gamma &< 0.06. 
\end{align*} \quad (20)
\]

These bounds are significantly larger than the values estimated above, especially those concerning the \( Z \)-interaction. If we take \( C, D \) to be in the upper range (see Eq. (11)), one finds \( BR(t \to c\gamma) \) and \( BR(t \to cZ) \) to be of the order of 1%. For our subsequent discussion we shall take this branching ratio as a guideline.

The Production of Single \( t \)-quarks at LEP II

If the anomalous vertices given in Eq. (10) are present, it is possible to produce single \( t \)-quarks in \( e^+ e^- \)-annihilation above about 180 GeV via the reaction \( e^+ e^- \to \gamma, Z \to \bar{t}t, \bar{t}c \). This possibility has also been discussed in Ref. [10] based on a different set of anomalous operators.

The total cross section for the single \( t \)-production is given by:

\[ \sigma_{\text{tot}} \left( e^+ e^- \to \bar{t}c, \bar{t} \bar{c} \right) = \sigma_\gamma + \sigma_Z + \sigma_{\text{int}} \quad (21) \]

with:

\[ \sigma_\gamma = \frac{N_C e^2}{16 \pi m_t^2} \left( |C_\gamma|^2 + |D_\gamma|^2 \right) c_\beta, \quad (22) \]

\[ \sigma_Z = \frac{N_C G_F m_Z^2}{16 \pi \sqrt{2} m_t^2} \left( 1 - 4 s_w^2 + 8 s_w^4 \right) \left( |C_Z|^2 + |D_Z|^2 \right) \frac{s^2}{(s - M_Z^2)^2 + \Gamma_Z^2 M_Z^2} c_\beta, \quad (23) \]

\[ \text{†While completing this work an additional preprint has appeared on the same issue [11].} \]

\[ \text{‡In the following formula for } \sigma_{\text{int}} \text{ the imaginary parts of the form factors are kept for completeness only. In calculating the cross section it was assumed that these imaginary parts vanish and that } C \text{ and } D \text{ are of the same size. These assumptions are made for definiteness only and should not influence the results qualitatively.} \]
\[ \sigma_{\text{int}} = \frac{N_{C}e}{16\pi\sqrt{2}m_{t}^{2}} \sqrt{\frac{G_{Z}M_{Z}^{2}}{2}} (1 - 4s_{w}^{2}) \frac{s}{(s - M_{Z}^{2})^{2} + \Gamma_{Z}^{2}M_{Z}^{2}} c_{\beta} \]
\[ \left[ 2(s - M_{Z}^{2}) (ReC_{\gamma}ReC_{Z} + ImC_{\gamma}ImC_{Z} + ReD_{\gamma}ReD_{Z} + ImD_{\gamma}ImD_{Z}) 
- 2\Gamma_{Z}M_{Z}(ImC_{\gamma}ReC_{Z} - ReC_{\gamma}ImC_{Z} + ImD_{\gamma}ReD_{Z} - ReD_{\gamma}ImD_{Z}) \right] . \quad (24) \]

The following abbreviations have been used:
\[ c_{\beta} = \beta_{+}^{2} \left( 4\beta_{\pm}^{2} - \frac{2}{3}\beta_{+}^{2} - 2\beta_{-}^{4} + 8\frac{m_{t}m_{c}}{s} \right) , \]
and
\[ \beta_{+}^{2} = 1 - \frac{(m_{t} + m_{c})^{2}}{s} , \quad \beta_{-}^{2} = 1 - \frac{(m_{t} - m_{c})^{2}}{s} , \quad \beta_{\pm}^{2} = 1 - \frac{m_{t}^{2} - m_{c}^{2}}{s} . \]

Numerically we find for \( BR_{\gamma} = BR_{Z} = 1\% \) at \( \sqrt{s} = 190\text{GeV} : \sigma_{\text{tot}} = 6.3 \cdot 10^{-2}\text{pb} \) \( (25) \)
and at \( \sqrt{s} = 200\text{GeV} : \sigma_{\text{tot}} = 13.2 \cdot 10^{-2}\text{pb} . \) \( (26) \)

In case of an integrated luminosity of 170 \( \text{pb}^{-1}/\text{year} \) one has 11 events or 22 events respectively. In view of the uncertainty of the vertex parameters \( C, D \) very detailed prediction cannot be made, but our result is encouraging. It seems not totally hopeless to look for \( t \)-quarks produced singly at LEP II. Of course, if this production is observed, it would be a clear indication towards a violation of the Standard Model \[12\].

**Single Production of t-quarks at HERA**

It is well-known that in ep-collisions at high energy studied at the experiments at HERA \( t \)-quarks cannot be produced at an observable rate if the Standard Model is valid \[13\]. In our case the reaction \( ec \to et \) can take place. Note that the c-quark is present inside the nucleon as part of the \( q\bar{q} \)-sea. The differential cross section is given by:

\[ \frac{d\sigma(ec \to et)}{dt} = \frac{1}{16\pi} \sum_{V,V'=\gamma,Z} g_{V}g_{V'}^{*} \frac{f_{V}}{\Lambda_{V}} \frac{f_{V'}}{\Lambda_{V'}} \frac{1}{D_{V}(t)D_{V'}(t)} \cdot \]
\[ \left\{ (C_{V}C_{V'}^{*} + D_{V}D_{V'}^{*}) (v_{V}v_{V'} + a_{V}a_{V'}) \left[ -(1 + \beta^{4}) t - \left( 1 + \beta^{2} \right) \frac{t^{2}}{s} \right] 
+ (C_{V}D_{V'}^{*} + D_{V}C_{V'}^{*}) (v_{V}a_{V'} + a_{V}v_{V'}) \left[ (1 - \beta^{4}) t + \left( 1 - \beta^{2} \right) \frac{t^{2}}{s} \right] \right\} . \quad (27) \]
For definiteness $\frac{f_{V}}{\Lambda_{V}}$ and $\frac{f_{V'}}{\Lambda_{V'}}$ are both fixed to $\frac{1}{m_{t}}$ in the following numerical examples. Further constants are $g_{\gamma} = (-i)Q_{e}$, where $Q_{e}$ is taken to be $-e$, and $g_{Z} = \left(\frac{i}{\sqrt{2}}\right)\sqrt{\frac{G_{F}M_{Z}^{2}}{\sqrt{2}}}$. The propagators in the t-channel are given by $D_{\gamma}(t) = t$ and $D_{Z}(t) = t - M_{Z}^{2}$ respectively. The constants $a_{V}$ and $v_{V}$ are the standard axialvector and vector coupling constants of $e^{-}$ or $e^{+}$ with the boson $V = \gamma, Z$ and finally $\beta^{2} = 1 - m_{t}^{2}/s$.

The dominant contribution comes from the $\gamma$-exchange which is given by:

$$
\sigma_{\gamma}(e c \rightarrow e t) dt = \frac{1}{8\pi m_{t}^{2}} \left( |C_{\gamma}|^{2} + |D_{\gamma}|^{2} \right) \left[ \frac{(1 + \beta^{4})}{t} - \frac{1 + \beta^{2}}{s} \right].
$$

(28)

It becomes obvious that the cross section gets the largest where $t$ is the smallest. Note that $|t|_{\text{min}}$ is proportional to $m_{e}^{2}$ in the limit $s \gg m_{t}^{2}$, namely

$$
|t|_{\text{min}} \simeq \frac{m_{e}^{2}(m_{t}^{2} - m_{e}^{2})^{2}}{s^{2}}.
$$

(29)

The dominant contribution to the production cross section thus comes from the phase space region near the threshold, where in the quark-lepton c.m.s. the incoming lepton looses all its momentum; the energy is transferred to the quark system in order to produce the heavy t-quark nearly at rest - in that particular frame. In the laboratory system both the t-quark and the lepton have no sizable transverse momentum, and the lepton will be lost in the forward direction. The t-quark will decay emitting a b-quark and a W-boson.

In order to calculate the cross section for $e p \rightarrow e t X$, it is necessary to integrate over the charm quark distribution function. In case of HERA, where the incoming lepton has an energy of 27.6 GeV, the c-quark needs to have a momentum of more than 277 GeV ($x > 0.338$) such that a t-quark can be produced. Thus the cross section is particularly large, if the c-quark distribution is as stiff as possible.

If one utilizes e.g. the charm distributions that Martin et al. have extracted from a global analysis [14] and assuming $BR_{\gamma} = BR_{Z} = 1\%$, the cross section is about $7 \cdot 10^{-4}pb$ which is too small to be observable at HERA. This conclusion remains the same even if the parameters $C$ and $D$ are allowed to saturate their bounds in Eq. (20) because these distributions fall off quickly for large $x$.

Alternatively we have considered charm distributions with a different large-$x$ behavior. Of particular interest is the “intrinsic-charm”-case [15]. The event rates have been calculated for the following two distribution functions that are based on the “intrinsic-charm” hypothesis [15, 16]:

$$
c_{1}(x) = \frac{1}{2} N_{5} x^{2} \left[ \frac{1}{3} (1 - x)(1 + 10x + x^{2}) + 2x(1 + x) \ln x \right],
$$

(30)
\[ c_2(x) = \frac{1}{210} N_5 x^8 \left[ 35 + 1155x - 1575x^2 - 11375x^3 - 2450x^4 + 490x^5 - 98x^6 + 14x^7 - x^8 + x \left( 1443 + 7161x + 5201x^2 \right) \right. \]
\[ \left. + \left\{ 840 + 5880x + 5880x^2 \right\} \ln x \right] . \] (31)

(Note: \( N_5 = 36 \) \((N_5 = 288028)\) have been used such that \( c(x) \) is normalized to 1%. The distribution function \( \bar{c}(x) \) is identical to \( c(x) \).) Again making use of the assumption \( BR_\gamma = BR_Z = 1\% \) one finds:

\[ c_1(x) : \quad \sigma_{tot} = 1.7 \cdot 10^{-2} pb , \] (32)

\[ c_2(x) : \quad \sigma_{tot} = 3.4 \cdot 10^{-2} pb . \] (33)

With an expected integrated luminosity of about 100 \( pb^{-1} \)/year one could therefore hope to observe a few events of this type for sufficient running time and favorable circumstances.

The topology of these events is such that it could not be explained within the Standard Model. If the \( W \) decays leptonically, one would observe a \( b \)-jet (or \( \bar{b} \)-jet) plus a high energy electron or muon.

It is interesting to note that the H1-Collaboration has recently reported the observation of several events with a high energy isolated lepton \(^{[17]}\). Three of these events in a data sample of 36 \( pb^{-1} \) show “kinematic properties atypical of Standard Model processes”. The production of single \( t \) (or \( \bar{t} \)) quarks could in principle generate this type of topology and kinematics. As a test of this hypothesis it would be mandatory to search for those events in which the \( W \) decays hadronically. For these events one expects to see three jets with jet momenta that should match to give the mass and momentum of the decaying top-quark.

Acknowledgement: We would like to thank Dr. S. Brodsky and Dr. A. Leike for useful discussions.

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