ELECTROWEAK DISCUSSION SECTION SUMMARY

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ABSTRACT

I summarize our work of electroweak discussion section held at the 14th International Workshop on Weak Interactions and Neutrinos. We discussed about a few physics topics related to electroweak interactions, including $B$ and $K$ physics and weak flavor physics.

1. INTRODUCTION

In the standard $SU(2) \times U(1)$ gauge theory of Glashow, Salam and Weinberg the fermion masses and hadronic flavor changing weak transitions have a somewhat less secure role, since they require a prior knowledge of the mass generation mechanism. The simplest possibility to give mass to the fermions in the theory makes use of Yukawa interactions involving the doublet Higgs field. These interactions give rise to the Cabibbo–Kobayashi–Maskawa (CKM) matrix: Quarks of different flavor are mixed in the charged weak currents by means of an unitary matrix $V$. However, both the electromagnetic current and the weak neutral current remain flavor diagonal. Second order weak processes such as mixing and CP–violation are even less secure theoretically, since they can be affected by both beyond the Standard Model virtual contributions, as well as new physics direct contributions. Our present understanding of CP–violation is based on the three–family Kobayashi–Maskawa model\textsuperscript{1}) of quarks, some of whose charged–current couplings have phases. Over the past decade, new data have allowed one to refine our knowledge about parameters of this matrix $V$.

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In Section 2, we introduce the quark mixing matrix, and a new set of quark mass matrices is proposed based on experimental mass hierarchy. In Section 3, we point out that directly measuring the invariant mass of final hadrons in $B$-meson semileptonic decays offers an alternative way to select $b \rightarrow u$ transitions that is in principle more efficient than selecting the upper end of the lepton energy spectrum. In Section 4, we briefly discuss about the role of vector mesons in rare kaon decays, and about probing supergravity models from the $b \rightarrow s\gamma$ decays.

2. FLAVOR DEMOCRACY AND FLAVOR GAUGE THEORY

It has turned out that for many fields of physics symmetry breaking is as fundamental a feature as symmetry itself. Whether one considers for instance the breakdown of rotational symmetry in a ferromagnet or the collapse of gauge symmetry in a superconductor, it is clear that our world would not be the same if these symmetries were respected at all temperatures. More fundamentally, it is the asymmetry of the electroweak vacuum that guarantees our very existence: in the absence of a spontaneous gauge symmetry breaking mechanism, quarks and pions would remain massless, and nuclear matter unbound. The hierarchical pattern of the quark masses and their mixing, nevertheless remains an outstanding issue of the electroweak theory. While a gauge interaction is characterized by its universal coupling constant, the Yukawa interactions have as many coupling constants as there are fields coupled to the Higgs boson. There is no apparent underlying principle which governs the hierarchy of the various Yukawa couplings, and, as a result, the standard model of strong and electroweak interactions can predict neither the quark (or lepton) masses nor their mixing.

The gauge invariant Yukawa Lagrangean involving quark fields is

$$\mathcal{L}_Y = - \sum_{i,j} \left( \bar{Q}'_{iL} \Gamma^{D}_{ij} d'_{jR} \phi + \bar{Q}'_{iL} \Gamma^{U}_{ij} u'_{jR} \tilde{\phi} + h.c. \right),$$

where the primed quark fields are eigenstate of the $SU(2) \times U(1)$ electroweak gauge interaction. The left-handed quarks form a doublet under the $SU(2)$ transformation, $\bar{Q}'_{L} = (\bar{u}'_{L}, \bar{d}'_{L})$, and the right-handed quarks are singlets. The indices $i$ and $j$ run over the number of fermion generations. The Yukawa coupling matrices $\Gamma^{U,D}$ are arbitrary and not necessarily diagonal. The Higgs field $\phi$ is parametrized in the unitary gauge as

$$\phi = \frac{1}{\sqrt{2}} \begin{pmatrix} 0 \\ v + H \end{pmatrix}. \quad (2)$$

After spontaneous symmetry breaking it acquires a non-vanishing vacuum expectation value
\( v \) which yields quark mass terms in the original Lagrangean

\[
\mathcal{L}_{\text{mass}} = -\frac{v}{\sqrt{2}} \sum_{i,j} (\bar{d}'_i L \Gamma^D_{ij} d'_j R + \bar{u}'_i L \Gamma^U_{ij} u'_j R + \text{h.c.}) .
\]  

(3)

The quark mass matrices are defined as

\[
M^{U,D}_{ij} \equiv \frac{v}{\sqrt{2}} \Gamma^{U,D}_{ij} .
\]  

(4)

These mass matrices \( M^{U,D} \) can be diagonalized with the help of unitary matrices, \( U^U_L \) and \( U^U_R \), and the electroweak eigenstates are transformed to physical mass eigenstates by the same unitary transformations,

\[
U^{U,D} M^{U,D} (U^{U,D})^\dagger = M^{U,D}_{\text{diag}} ,
\]  

(5)

The kinetic and neutral current terms of electroweak gauge interactions remain invariant under the unitary transformation (5). However, the charged current terms display non-diagonal couplings when written in terms of the physical quarks. Defining the CKM matrix as

\[
V = U^U U^{D\dagger} ,
\]  

(6)

the charged current Lagrangean becomes

\[
\mathcal{L}_{\text{c.c.}} \sim \sum_{i,j,k} \left[ \bar{u}_i L \gamma^\mu V_{ij} d_{jL} W^+_{\mu} + \bar{d}_i L \gamma^\mu V_{ij}^\dagger u_{jL} W^-_{\mu} \right] .
\]  

(7)

In the past, many attempts at unifying the gauge interactions of the standard model (SM) have been made, within the framework of the Grand Unified Theories (GUT). These theories yield the unification energy \( E_{\text{GUT}} \geq 10^{16} \text{ GeV} \), i.e. the energy where the SM gauge couplings are to coincide:

\[
\frac{5}{3} \alpha_1 = \alpha_2 = \alpha_3 .
\]  

(8)

Here, \( \alpha_j = g_j^2 / 4\pi \), and \( g_1, g_2, g_3 \) are the gauge couplings of \( U(1)_Y \), \( SU(2)_L \), \( SU(3)_C \), respectively. For the condition (8) to be satisfied at a single point \( \mu (= E_{\text{GUT}}) \) exactly, supersymmetric theories (SUSY) were introduced to replace SM at \( \mu \approx M_{\text{SUSY}} \approx 1 \text{ TeV} \). This changed the slopes of the functions \( \alpha_j = \alpha_j(\mu) \) at \( \mu \geq M_{\text{SUSY}} \), and for certain
values of parameters of SUSY the three lines met at a single point, as required by (8). However, such an approach has several deficiencies. Due to large lower bound for the proton decay time \( \tau_{\text{proton}} \geq 5.5 \times 10^{32} \text{ yr} \), the unification energy is exceedingly large \( E_{\text{GUT}} \geq 10^{16} \text{ GeV} \). This would imply that there is a large desert between \( M_{\text{SUSY}} \) and \( E_{\text{GUT}} \). Furthermore, although the initial aim of the attempts leading to SUSY and GUT was to reduce the number of independent parameters (gauge couplings), the overall number of free parameters was in fact substantially increased through the inclusion of SUSY degrees of freedom (sfermions, additional Higgses, etc.).

We believe that it is more reasonable to attempt first to reduce the number of degrees of freedom in the sector of Higgses and Higgs-fermion interactions (i.e. Yukawa sector), since this sector seems to be more problematic than the gauge boson sector. Furthermore, we believe that any such attempt should be required to lead to an overall reduction of the seemingly independent degrees of freedom, unlike the GUT–SUSY approach. A flavor gauge theory (FGT)\(^2\) is a theory having universal strengths of the Yukawa (running) couplings at energies \( E \geq \Lambda_{\text{FGT}} \). The symmetry responsible for this reduction of the number of parameters would be “flavor democracy” (FD), valid possibly in certain separate sectors of fermions. Although FGT could imply that the Higgs(es) of SM are elementary at low energies, it appears natural that FGT is without Higgs and that the Higgs particles of SM are bound states of fermion pairs \( \langle \bar{f}f \rangle \), created through dynamical symmetry breaking (DSB) in a transition interval \([E_{\text{trans.}}, \Lambda_{\text{FGT}}]\) between SM and FGT. The idea of FD and the deviation from the exact FD at the low energy \( E \sim 1 \text{ GeV} \) has been investigated by several authors\(^3\). On the other hand, here we investigate the deviation at higher energies and the possible connection to FGT.

We can illustrate these concepts with a simple scheme of an FGT. Assume that at energies \( E \geq \Lambda_{\text{FGT}} \) we have no Higgses, but a new gauge boson \( B_\mu \), with a mass \( M_B (> \Lambda_{\text{FGT}}) \), i.e. the symmetry group of the extended gauge theory is \( G_{\text{SM}} \times G_{\text{FGT}} \). The SM–part \( G_{\text{SM}} \) is \( SU(3)_C \times SU(2)_L \times U(1)_Y \) without Higgses, and hence with (as yet) massless gauge bosons and fermions at these energies. The FGT–part of Lagrangian in the fermionic sector is written schematically as

\[
\mathcal{L}^\text{FGT}_{G-\bar{f}} = -g \Psi \gamma^\mu B_\mu \Psi \quad \text{(for } E > \Lambda_{\text{FGT}}) ,
\]  

where \( \Psi \) is the column of all fermions and \( B_\mu = B^j_\mu T_j \), \( T_j \)'s being the generator matrices of the new symmetry group \( G_{\text{FGT}} \). Furthermore, we assume that the \( T_j \)'s corresponding to the electrically neutral \( B^j_\mu \)'s do not mix flavors (i.e. no FCNC at tree level) and are proportional to identity matrices in the flavor space. We will show in the following lines that the FGT Lagrangian (9) can imply the creation of SM–Higgs particles through DSB (i.e. condensation of fermion pairs), as well as the creation of Yukawa couplings with a flavor democracy.
The effective current–current interaction, corresponding to the exchange of neutral gauge boson \( B \) at low energies \( E (\Lambda_{\text{FGT}} \leq E < M_B) \), is

\[
L_{\text{FGT}}^{4f} \approx -\frac{g^2}{2M_B^2} \sum_{i,j} (\bar{f}_i \gamma^\mu f_i)(\bar{f}_j \gamma_\mu f_j) \quad (\text{for } \Lambda_{\text{FGT}} \leq E < M_B).
\]

(10)

Since we are interested in the possibility of Yukawa interactions of SM being contained in (10), and since such interactions connect left–handed to right–handed fermions, we have to deal only with the left–to–right (and right–to–left) part of (10). By applying Fierz transformation\(^4\) to this part, we end up with the 4–fermion interactions without \( \gamma_\mu \)’s

\[
L_{\text{FGT}}^{4f} \approx \frac{2g^2}{M_B^2} \sum_{i,j} (\bar{f}_i L f_j R)(\bar{f}_j R f_i L) \quad (\text{for } \Lambda_{\text{FGT}} \leq E < M_B).
\]

(11)

These interactions can be rewritten in a mathematically equivalent (Yukawa) form with auxiliary (i.e. as yet non–dynamical) scalar fields. One possibility is to introduce only one isodoublet auxiliary scalar \( H \) with (as yet arbitrary) bare mass \( M_H \), by employing a familiar mathematical trick\(^5\)

\[
\mathcal{L} \approx -M_H \sqrt{2g} \sum_{i,j=1}^3 \left[ (\bar{\psi}_i^q L H) u^q_{jR} + (\bar{\psi}_i^l L H^\dagger) u^l_{jR} + h.c. \right] \\
+ \left[ (\bar{\psi}_i^q L H) d^q_{jR} + (\bar{\psi}_i^l L H) d^l_{jR} + h.c. \right] \\
- M_H^2 (H^\dagger H),
\]

(12a)

where \( H = \begin{pmatrix} H^\dagger \\ H^0 \end{pmatrix}, \quad \tilde{H} = i\tau_2 H^* \),

\[
\psi_i^q = \begin{pmatrix} u^q_i \\ d^q_i \end{pmatrix}, \quad \psi_i^l = \begin{pmatrix} u^l_i \\ d^l_i \end{pmatrix},
\]

(12b)

\[
u_e^i, \quad \nu^i \mu, \quad \nu^i \tau, \quad d^q_i = \begin{pmatrix} d \\ s \\ b \end{pmatrix}, \quad d^l_i = \begin{pmatrix} \mu^- \\ \tau^- \end{pmatrix}.
\]

Another possibility is to introduce two auxiliary scalar isodoublets \( H^{(U)}, H^{(D)} \), with (as
yet) arbitrary bare masses $M_H^{(U)}$, $M_H^{(D)}$, and express (11) in the 2-Higgs ‘Yukawa’ form

$$\mathcal{L} \approx - M_H^{(U)} \sqrt{2} g \frac{\sqrt{2} g}{M_B} \sum_{i,j=1}^{3} \left[ \left( \tilde{\psi}_{iL}^q \tilde{H}^{(U)} \right) u_j^q + \left( \tilde{\psi}_{iL}^l \tilde{H}^{(U)} \right) u_j^l + h.c. \right]$$

$$- M_H^{(D)} \sqrt{2} g \frac{\sqrt{2} g}{M_B} \sum_{i,j=1}^{3} \left[ \left( \tilde{\psi}_{iL}^q \tilde{H}^{(D)} \right) d_j^q + \left( \tilde{\psi}_{iL}^l \tilde{H}^{(D)} \right) d_j^l + h.c. \right]$$

$$- M_H^{(U)}^2 \left( H^{(U)} \right) \left( H^{(U)} \right)^\dagger - M_H^{(D)}^2 \left( H^{(D)} \right) \left( H^{(D)} \right)^\dagger .$$

(13)

As derived, the Yukawa couplings (12a) and (13) are valid in the FGT–low energy region ($\Lambda_{\text{FGT}} \leq E < M_B$), having non–dynamical scalar fields and being mathematically equivalent to (11). However, with decreasing the energy, the scalars in (12) and (13) can be shown to obtain vacuum expectation values (VEV’s) and kinetic terms through quantum effects. The neutral components of the physical Higgs doublets created through DSB in this way are then condensates of fermion pairs $\langle \bar{f} f \rangle$. (Note that this scheme of DSB is just one of many possible scenarios for an FGT–SM transition. However, our results will not depend on any such specific scenario, but rather on the assumption of FD near the transition energies, as expressed in (12) and (13).) This would then lead us to the minimal SM (with one Higgs) in the case (12) and to SM with two Higgses (type II) in the case (13). Hence, although (12) and (13) are mathematically equivalent, they lead to two physically different theories.

Note that (12) implies that the minimal SM, if it is to be replaced by FGT at high energies, should show up a trend of the Yukawa coupling matrix (or equivalently: of the mass matrix in the flavor basis) toward a complete flavor democracy for all fermions, with a common overall factor, as the energy is increased within the minimal SM toward transition region denoted as $E_0 (\equiv E_{\text{trans.}})$

$$M^{(U)} \quad \text{and} \quad M^{(D)} \to \frac{1}{3} m_t^0 \left( \begin{array}{cc} N_{FD}^q & 0 \\ 0 & N_{FD}^l \end{array} \right) \quad \text{as} \quad E \uparrow E_0 , \quad (14a)$$

where $m_t^0 = m_t (\mu = E_0)$ and $N_{FD}$ is the $3 \times 3$ flavor–democratic matrix

$$N_{FD}^f = \frac{1}{3} \left( \begin{array}{ccc} 1 & 1 & 1 \\ 1 & 1 & 1 \\ 1 & 1 & 1 \end{array} \right) , \quad (14b)$$

with the superscript $f = q$ for the quark sector and $f = l$ for the leptonic sector. On the other hand, if SM with two Higgses (type II) is to experience such a transition, then (13)
implies separate trends toward FD for the up-type and down-type fermions

\[ M^{(U)} (M^{(D)}) \rightarrow \frac{1}{3} m_t^0 (m_b^0) \left( \begin{array}{cc} N_q^{FD} & 0 \\ 0 & N_{FD}^l \end{array} \right) \text{ as } E \uparrow E_0 , \quad (15) \]

where \( m_t^0 \) and \( m_b^0 \) can in general be different. Note that \( N_{FD} \), when written in the diagonal form in the mass basis, acquires the form

\[ N_{FD}^{mass \ basis} = 3 \left( \begin{array}{ccc} 0 & 0 & 0 \\ 0 & 0 & 0 \\ 0 & 0 & 1 \end{array} \right) . \quad (16) \]

Hence, FD (and FGT) implies in the mass basis as \( E \) increases to \( E_0 \),

\[ \frac{m_u}{m_t}, \frac{m_c}{m_t}, \frac{m_{\nu_e}}{m_{\nu_e}}, \frac{m_{\nu_\mu}}{m_{\nu_\mu}} \rightarrow 0 , \quad (17a) \]
\[ \frac{m_d}{m_b}, \frac{m_s}{m_b}, \frac{m_{\nu_s}}{m_{\nu_s}}, \frac{m_{\nu_\tau}}{m_{\nu_\tau}} \rightarrow 0 , \quad (17b) \]
\[ \frac{m_{\nu_\mu}}{m_t}, \frac{m_{\nu_\tau}}{m_b} \rightarrow 1 , \quad (17c) \]

and in the case of the minimal SM in addition

\[ \frac{m_b}{m_t}, \frac{m_{\nu_\tau}}{m_{\nu_\tau}} \rightarrow 1 . \quad (18) \]

In Ref. 2, we have shown that the minimal SM does not show the required trend toward FD, but that SM with two Higgs doublets (type II) does. For more details, see Ref. 2.

3. HADRONIC INVARIANT MASS DISTRIBUTION ON SEMILEPTONIC \( B \)-DECAY

Semileptonic decays of \( B \)-mesons continue to attract great experimental and theoretical interest, because of the light they may throw both on heavy quark decay mechanisms and on KM matrix elements. Theoretical approaches have been made both from the inclusive point of view, as in the independent-quark decay model of Altarelli et al.\(^7\), and also by calculating and summing exclusive channels as with Wirbel et al.\(^8\) and others. Here we adopt the inclusive approach.

In the independent-quark decay model (or “spectator model”) of Ref. 7, the initial \( B \)-meson is represented by a pair of quarks \( b + \bar{q} \) (\( \bar{q} = \bar{u} \) or \( \bar{d} \)) with a distribution of Fermi momentum \( p = p(\bar{q}) = -p(b) \) in the \( B \) rest-frame. Then \( b \) decays as a free quark via
$b \to q'\ell\nu$ (with $q' = c$ or $u$); $\bar{q}$ remains an independent spectator. The spectator mass is assigned a value of order $m_{\text{sp}} \simeq 0.1$ GeV (the precise value is not critical). Kinematical constraints then require the effective mass of the decaying $b$-quark to depend on the Fermi momentum $p$: 

$$m_B^2 = m_B^2 + m_{\text{sp}}^2 - 2m_B(m_{\text{sp}}^2 + p^2)^{1/2},$$  \hspace{1cm} (19)$$

where $p = |\mathbf{p}|$. The primary role of the Fermi momentum is to reflect the initial $B(b\bar{q})$ wave function, through the impulse approximation; however, in the spectator picture, it also subsumes some effects of final state interactions. For example, large values of $p$ would contribute toward broad distributions in the final hadronic invariant mass $m(\bar{q}q')$; if final-state interactions enhance a relatively low-mass region, their effects could appear as an enhancement of the effective Fermi momentum distribution at small $p$.

The lepton spectrum in $b \to c\ell\nu$ decays has the form

$$\frac{d\Gamma(b \to c\ell\nu)}{dx_\ell} = \frac{d\Gamma^0}{dx_\ell} \left[ 1 - \frac{2\alpha_s}{3\pi} G(x_\ell, \epsilon) \right]$$ \hspace{1cm} (20)$$

where $x_\ell = 2(\ell \cdot b)/m_b^2 = 2E_\ell/m_b$ in the $b$ rest-frame, $\Gamma^0$ is the standard lowest-order electroweak calculation, $\alpha_s$ is the QCD coupling constant and $\epsilon = m_c/m_b$. The factor in square brackets is the lowest-order QCD correction from virtual and real gluon emission and $G(x_\ell, \epsilon)$ is a complicated function defined in Ref. 7, including a Sudakov resummation of double logarithms to ensure a finite result at the end-point $x_\ell(\text{max}) \equiv x_M = 1 - \epsilon^2$. $G$ depends little on $x_\ell$ except near $x_M$; for other values of $x_\ell$, $G$ is close to its spectrum-averaged value $g(\epsilon)$, defined by the integrated width

$$\Gamma(b \to c\ell\nu) = \Gamma^0(b \to c\ell\nu) \left[ 1 - \frac{2\alpha_s}{3\pi} g(\epsilon) \right].$$ \hspace{1cm} (21)$$

There is a useful empirical approximation\(^9\) accurate within 0.2%,

$$g(\epsilon) \simeq (\pi^2 - 31/4)(1 - \epsilon)^2 + 3/2.$$ \hspace{1cm} (22)$$

For completeness we note that

$$\Gamma^0(b \to c\ell\nu) = \frac{G_F^2 m_b^5}{192 \pi^3} |V_{cb}|^2 (1 - 8\epsilon^2 + 8\epsilon^6 - \epsilon^8 - 24\epsilon^4 \ln \epsilon),$$ \hspace{1cm} (23)$$

$$\frac{d\Gamma^0}{dx_\ell}(b \to c\ell\nu) = \frac{G_F^2 m_b^5}{96 \pi^3} |V_{cb}|^2 x_\ell^2 \frac{(1 - x_\ell^2 - x_\ell)^2}{(1 - x_\ell)^3} \left[ (1-x_\ell)(3-2x_\ell) + \epsilon^2(3-x_\ell) \right],$$ \hspace{1cm} (24)$$

where $V$ is the CKM matrix and lepton mass is ignored. For each value of the Fermi momentum $\mathbf{p}$, we calculate $d\Gamma/dE_\ell$ in the $b$ rest-frame (isotropic angular distribution) and
boost to the \( B \) rest-frame folded with Fermi motion \( p \). For comparison with data at the \( \Upsilon(4S) \) resonance, where \( B \) is produced not precisely at rest, a further boost to the lab frame is made with appropriate angular integrations.

The CKM matrix element \( V_{ub} \) characterizing \( b \to u \) quark transitions plays an important role in the description of CP violation within the three-family Standard Model, but is still not accurately known. The most direct way to determine this parameter is through the study of \( B \) meson semileptonic decays; recent results from the CLEO\(^{10}\) and ARGUS\(^{11}\) data on the end-point region of the lepton spectrum have established that \( V_{ub} \) is indeed non-zero and have given an approximate value for its modulus. The central problem in the extraction of \( V_{ub} \) is the separation of \( b \to u \) events from the dominant \( b \to c \) events. In semi-leptonic \( B \)-meson decays, the usual approach is to study the upper end of the charged lepton spectrum, since the end-point region

\[
E_\ell > (m_B^2 - m_Y^2) / (2m_B)
\]

in the CM frame is inaccessible to \( b \to c \) transitions and therefore selects purely \( b \to u \). However, only about 20\% of \( b \to u \) transitions actually lie in the region of Eq. (25); it is therefore not a very efficient way to select them. In situations of physical interest, the situation is even somewhat worse. For example, in \( \Upsilon(4S) \to BB \) decay, each \( B \) meson has a small velocity in the \( \Upsilon \) rest frame; the magnitude \( \beta \) of this velocity is known, but its direction is not. In this frame, which is the laboratory frame when \( \Upsilon \) is produced at a symmetric \( e^+e^- \) collider, the \( b \to u \) selection region based on Eq. (25) becomes

\[
E_\ell > \gamma (1 + \beta) (m_B^2 - m_Y^2) / (2m_B),
\]

for the cases \( \ell = e \) or \( \mu \). Here \( \gamma = (1 - \beta^2)^{-1/2} = m_\Upsilon / (2m_B) \) and we neglect the lepton mass. (At an asymmetric collider, where \( e^+ \) and \( e^- \) beams have different energies, it will be necessary to boost lepton momenta from the laboratory frame to the \( \Upsilon \) rest frame before applying this cut.) Equation (26) accepts an even smaller percentage of \( b \to u \) decays than Eq. (25), about 10\% in fact.

The essential physical idea behind Eqs. (25)-(26) is that \( b \to c \) transitions leave at least one charm quark in the final state; hence for a general semileptonic decay \( B \to \ell + \nu + X \) the invariant mass \( m_X \) of the final hadrons exceeds \( m_D \) and this implies a kinematic bound on \( E_\ell \). In Ref. 12, we give this old idea a new twist. We first observe that there is no unique connection between \( m_X \) and \( E_\ell \), due to the presence of the neutrino, so the bound on \( E_\ell \) is not an efficient way of exploiting the bound on \( m_X \). We then observe a more efficient way to exploit the latter bound is to measure \( m_X \) itself and to select \( b \to u \) transitions by requiring

\[
m_X < m_D
\]

instead of Eqs. (25)-(26). This condition is of course frame-independent.
The final hadronic invariant mass distribution depends both on the $c$-quark energy distribution $d\Gamma(b \to c\ell\nu)/dE_c$ and on the Fermi momentum distribution $\phi(p)$ which is normalized to $\int_0^\infty dp\,\phi(p) = 1$ and is sometimes approximated by a gaussian form
\[
\phi(p) = 4p^2(p_F^3\sqrt{\pi})^{-1}\exp(-p^2/p_F^2)
\]
where $p_F$ is a characteristic parameter describing Fermi motion inside $B$. The lowest-order contribution to the $c$-quark energy distribution is given by
\[
d\Gamma^0(b \to c\ell\nu)/dx_c = \frac{G_F^2m_b^5}{96\pi^3}|V_{cb}|^2(x_c^2 - 4\epsilon^2)^{1/2}[3x_c(3 - 2x_c) + \epsilon^2(3x_c - 4)],
\]
where $x_c = 2(c\cdot b)/m_b^2 = 2E_c/m_b$ in the $b$ rest-frame, with kinematical range $2\epsilon \leq x_c \leq 1 + \epsilon^2$. When QCD radiative corrections are included, the real and virtual gluon contributions must be subject to resolution smearing so that their singular parts will cancel; this we approximate by absorbing real soft gluons into the effective final $c$-quark and correcting $d\Gamma^0/dx_c$ by the factor $g(\epsilon)$:
\[
\frac{d\Gamma(b \to c\ell\nu)}{dx_c} \approx \frac{d\Gamma^0(b \to c\ell\nu)}{dx_c} \left[1 - \frac{2\alpha_s}{3\pi}g(\epsilon)\right].
\]
For each value of the Fermi momentum $p$ we calculate $d\Gamma/dE_c$ in the $b$ rest-frame (isotropic angular distribution here) and fold it with the spectator energy-momentum vector to form the distribution $d\Gamma/dm_X$ with respect to the invariant mass $m_X$ of the final charmed hadronic system,
\[
m_X^2 = (E_c + E_{sp})^2 - (p_c + p_{sp})^2.
\]
The spectator energy and momentum in the $b$ rest-frame are
\[
E_{sp} = \left[(p^2 + m_b^2)^{1/2}(p^2 + m_{sp}^2)^{1/2} + p^2\right]/m_b,
\]
\[
p_{sp} = \left[(p^2 + m_b^2)^{1/2} + (p^2 + m_{sp}^2)^{1/2}\right]p/m_b,
\]
and $m_b$ is everywhere defined by Eq. (19). The maximum and minimum values of $m_X^2$ for given $p$ are
\[
m_X^2(\text{max}) = m_c^2 + m_{sp}^2 + m_b(E_{sp} + p_{sp}) + m_c^2(E_{sp} - p_{sp})/m_b,
\]
\[
m_X^2(\text{min}) = \begin{cases} (m_c + m_{sp})^2, & \text{if } (m_b^2 - m_c^2)E_{sp} \geq (m_b^2 + m_c^2)p_{sp}, \\ m_c^2 + m_{sp}^2 + m_b(E_{sp} - p_{sp}) + m_c^2(E_{sp} + p_{sp})/m_b, & \text{otherwise.} \end{cases}
\]
These relations show explicitly that small $p$ results in $m_X$ values close to $(m_c + m_{sp})$. The
upper limit on $p$ for the decay to be possible (from Eq. (19) with $m_b > m_c$) is

$$p < \lambda^{1/2}(m_B^2, m_c^2, m_{sp}^2)/(2m_B),$$

(34)

where $\lambda(a, b, c) = a^2 + b^2 + c^2 - 2ab - 2bc - 2ac$.

For $b \to u$ transitions the effect of individual resonances in $X$ quickly disappears above the $\pi$ and $\rho$ region, and multiparticle jet-like continuum final states should give the dominant contributions$^{13}$; this makes it reliable to calculate the $m_X$ distribution using the modified spectator decay model. Figure 1 gives the hadronic invariant mass distribution from $B \to \ell \nu X$ semileptonic decays, showing that more than 90% of $b \to u$ decays lie within the region selected by Eq. (27). In this illustration we use $m_B = 5.273$ GeV, $E_B = m_{\Upsilon}/2 = 5.29$ GeV, $m_c = 1.6$ GeV, $m_u = 0.1$ GeV, $p_F = 0.3$ GeV, including QCD corrections up to order $\alpha\alpha_s$ according to Ref. 14. The figure is normalized for simplicity to the case $|V_{ub}/V_{cb}| = 1$.

In order to exploit Eq. (27) instead, it is desirable to isolate uniquely the products of a single $B$ meson decay; to achieve this it is generally necessary to reconstruct both $B$ decays in a given event. One of these decays can be semileptonic, since kinematic constraints can often determine the missing neutrino four-momentum well enough to reconstruct a peak at zero in the invariant square of this four-momentum. Double-semileptonic decay events will not generally reconstruct uniquely, however. To study semileptonic channels, we are therefore concerned with those events (about 30% of the total) where one $B$ decays hadronically, one semileptonically with $\ell = e$ or $\mu$. Of order 1% of these have $b \to u \ell \nu$ semileptonic transitions (because of the very small ratio $|V_{ub}/V_{cb}| \approx 0.1$). About 10% of the latter satisfy the criterion $E_{\ell} > 2.5$ GeV of Eq. (26). With present data, it appears to be possible to reconstruct a few percent of such events, but perhaps only about one percent without ambiguity. The arithmetic of collecting reconstructed $b \to u \ell \nu$ events with the criterion of Eq. (26) is then approximately as follows:

$$10^6 B\bar{B} \text{ events} \to 3 \times 10^5 \text{ events with one } (b \to q\ell\nu)$$

$$\to 3 \times 10^3 \text{ events with one } (b \to u\ell\nu)$$

$$\to 3 \times 10^2 \text{ events with } (b \to u\ell\nu, E_{\ell} > 2.5 \text{ GeV})$$

$$\to 3 \text{ reconstructed } (b \to u\ell\nu, E_{\ell} > 2.5 \text{ GeV}).$$

(35)

The numbers in Eq. (35) are at first sight discouraging to hopes of an extensive study of $b \to u \ell \nu$ transitions.

Figure 1. Hadronic invariant mass distribution in $B \to \ell \nu X$ semileptonic decays.
In contrast, using the criterion of Eq. (27) instead offers an order of magnitude more events:

\[ 10^6 \bar{B}B \text{ events} \rightarrow 3 \times 10^3 \text{ events with } (b \rightarrow u\ell\nu, \ m_X < m_D) \]
\[ \rightarrow 30 \text{ reconstructed } (b \rightarrow u\ell\nu, \ m_X < m_D). \]  

(36)

This is more encouraging. Also, the introduction of microvertex detectors (ARGUS has recently added one) can be expected to improve the success rate for reconstructing \( \bar{B}B \) events and to reduce the ambiguous cases. For symmetrical colliders, microvertex detectors will assist in identifying charm decay vertices and the charged tracks associated with them. For projected future asymmetrical colliders, where the two beams have different energies and the produced \( \Upsilon(4S) \) is not at rest in the laboratory, the bottom decay vertices too will often be identifiable, with still greater advantage to the analysis. These future “\( b \)-factories” will also deliver much larger numbers of events; samples of order \( 10^8 \) or more \( \bar{B}B \) events are foreseen, in which case the numbers in Eq. (36) can be scaled up by at least two orders of magnitude.

We note that there is a question of bias. Some classes of final states (e.g. those with low multiplicity, few neutrals) may be more susceptible to a full and unambiguous reconstruction. Hence an analysis that requires this reconstruction may be biased. However the use of topological information from microvertex detectors should tend to reduce the bias, since vertex resolvability depends largely on the proper time of the decay and its orientation relative to the initial momentum (that are independent of the decay mode). Also such a bias can be allowed for in the analysis, via suitable modeling. Finally there may be a background from continuum events that accidentally fake the \( \Upsilon \) events of interest. This can be measured directly at energies close to the resonance.

4. TOPICS ON \( K \) AND \( B \) PHYSICS

4.1 Role of vector mesons in rare kaon decays\(^{15}\)

The role of vector mesons in \( K_L \rightarrow \pi^0\gamma\gamma \) and \( K_L \rightarrow \pi^0e^+e^- \) has been a controversial subject\(^{16,17}\). Chiral perturbation theory\(^{16}\) and the pion rescattering model\(^{18}\) predict that the pion loop gives a dominant contribution to \( K_L \rightarrow \pi^0\gamma\gamma \) with branching ratio around \( 7 \times 10^{-7} \), and the two photon spectrum has a peak near \( m_{\gamma\gamma} \approx 300 \text{ MeV} \) in both models. Also, the low energy photon pair is almost negligible. On the other hand, there can be a large enhancement in the low \( m_{\gamma\gamma} \) region\(^{19,20}\), if vector mesons come into play in this decay mode and one uses the naive nonet symmetry in \( K_2 \rightarrow P \) (with \( P = \pi^0, \eta_8, \eta_0 \)). In this case, the branching ratio is \( (1 \sim 3) \times 10^{-6} \). The recent measurement\(^{21}\) of \( K_L \rightarrow \pi^0\gamma\gamma \) from CERN is rather puzzling. The two photon spectrum seems consistent with the predictions of the chiral perturbation theory and the pion rescattering model. However, the branching ratio for \( m_{\gamma\gamma} \geq 280 \text{ MeV} \) is larger than those predictions by a factor of \( 3 \sim 4 \). In this rather confusing situation, it would be nice to have a systematic calculation for \( K_L \rightarrow \pi^0\gamma\gamma \) as well as for other processes where vector meson contributions can be potentially important.
In Ref. 15, the author gives self-consistent calculations for $K \to \pi\pi$, $K \to \pi l^+l^-$, $K_L \to \gamma\gamma$, $K_L \to \gamma l^+l^-$ and $K_L \to \pi^0\gamma\gamma$ in the hidden symmetry scheme in conjunction with the nonleptonic Hamiltonian used by J.J. Sakurai, J.A. Cronin and other groups. In his approach, the usual $O(p^4)$ and $O(p^6)$ terms arise from vector meson exchange between the chiral mesons, and from the Wess–Zumino anomaly term including vector mesons. From the work of Ref. 20, it is known that the $O(p^6)$ terms through $\rho^0$ and $\omega$ exchange in $\gamma\gamma \to \pi^0\pi^0$ unitarize the chiral loop amplitude in an effective way up to $m_{\gamma\gamma} \sim 1$ GeV, and controls the high energy behavior of the chiral amplitudes. Once we make this assumption, things get much simplified and we anticipate we do not lose any important physics information. The major advantage of our model lies in the number of unknown parameters to be determined by the experimental data. In fact, we do not have any unknown parameters at the level of strong and electromagnetic interactions other than the meson masses and their decay constants, *once* we adopt the notion of vector meson dominance in the normal sector. We first consider the structure of the left-handed currents.

We find that there exist anomalous left–handed currents arising from the Wess–Zumino anomaly and the intrinsic parity violating interactions involving vector mesons, and they give important contributions to $K_L \to \gamma l^+l^-$ and $K_L \to \pi^0\gamma\gamma$ through generating weak $V\pi\gamma$ vertices. The nonleptonic weak decays of kaons is described by the effective weak Lagrangian of current–current interactions. We make connections with the calculations in terms of quark fields with QCD corrections, and find that there is an additional term in the effective weak Lagrangian, which is proportional to $\text{Tr} (j_{\mu L})$ where ‘Tr’ is taken over the $U(3)_f$ or $SU(3)_f$ indices. This additional term cannot be thrown away as usually done once the masses and their decay constants, we find there exist anomalous left–handed currents arising from the Wess–Zumino anomaly and the intrinsic parity violating interactions involving vector mesons, and they give important contributions to $K_L \to \gamma l^+l^-$ and $K_L \to \pi^0\gamma\gamma$ through generating weak $V\pi\gamma$ vertices. The nonleptonic weak decays of kaons is described by the effective weak Lagrangian of current–current interactions. We make connections with the calculations in terms of quark fields with QCD corrections, and find that there is an additional term in the effective weak Lagrangian, which is proportional to $\text{Tr} (j_{\mu L})$ where ‘Tr’ is taken over the $U(3)_f$ or $SU(3)_f$ indices. This additional term cannot be thrown away as usually done in the case of $SU(3)_L \times SU(3)_R$, since the anomalous left–handed currents we consider here have nonvanishing trace even in that case. The coefficient of this new term measures the contributions of penguin operators of current–current types to the nonleptonic kaon decays.

At this stage, we will have four parameters, $C_8^{(1/2)}, C_27^{(1/2)}, C_27^{(3/2)}$ and $\delta_p$ with one constraint $C_27^{(3/2)} = 5C_27^{(1/2)}$. C’s characterize the strengths of the $(8_L, 1_R)_{\Delta I=1/2}, (27_L, 1_R)_{\Delta I=1/2}$ and $(27_L, 1_R)_{\Delta I=3/2}$ pieces of the weak Hamiltonian, respectively. Deviation of $\delta_p$ from 1 measures the contributions of penguin operators. We fix C’s from $K \to \pi\pi$. $\delta_p$ contributes to $K_L \to \pi^0\gamma\gamma$ in the $SU(3)_L \times SU(3)_R$ case, and both to $K_L \to \pi^0\gamma\gamma$ and to $K_L \to \gamma l^+l^-$ in the $U(3)_L \times U(3)_R$ case.

Next we assume nonet symmetry in the vector meson sector. Before studying the effects of $\delta_p$ on $K_L \to \gamma l^+l^-$, we analyze $K \to \pi l^+l^-$ using the suitable value of $\delta_p$. For the decay mode $K \to \pi l^+l^-$, we need two more operators, $Q_7V \equiv (\bar{s}d)_{V-A}(ll)V$ and $Q_7A \equiv (\bar{s}d)_{V-A}(ll)\bar{A}$, arising from the electromagnetic penguin diagram, the $Z^0$ penguin diagram and the box diagram with two internal $W$'s. The real part of the Wilson coefficients of these operators cannot be calculated reliably. Therefore, we introduce another parameter $C_7$ as the coefficient of the new operator $Q_7V$, ignoring the operator $Q_7A$. We fix $C_7$ from the best fit to the decay mode $K^+ \to \pi^+e^+e^-$. There is a twofold ambiguity in $C_7$, and we can predict for $K^+ \to \pi^+\mu^+\mu^-$, $K_S \to \pi^0e^+e^-$ and $K_S \to \pi^0\mu^+\mu^-$. This
model predicts the decay rate for \( K_S \rightarrow \pi^0 e^+ e^- \) process to be comparable to the decay rate for \( K^+ \rightarrow \pi^+ e^+ e^- \). The \( K_S \rightarrow \pi^0 e^+ e^- \) process contributes to the indirectly CP-violating \( K_L \rightarrow \pi^0 e^+ e^- \) process through the mixing between \( K_L \) and \( K_S \). We predict that the branching ratio of the indirectly CP-violating \( K_L \rightarrow \pi^0 e^+ e^- \) process is about \( 1.4 \) or \( 2.7 \times 10^{-10} \), which is substantially larger than other previous calculations. Another process in which vector mesons are important is \( K_L \rightarrow \gamma e^+ e^- \). This has been recently remeasured and the form factor shows a clear deviation from the \( \rho \) form factor\(^{28, 29} \). Then we make an analysis of \( K_L \rightarrow \gamma \gamma \) and \( K_L \rightarrow \gamma l^+ l^- \). In \( K_L \rightarrow \gamma l^+ l^- \), we have both weak \( VV \) and weak \( V \pi \gamma \) vertices, the latter of which was not considered in the earlier analysis. Next we introduce one more parameter \( \delta_n \) characterizing the possible deviation of \( a(K_2 \eta_0) \) from its naive value obtained from the effective weak Lagrangian. This takes care of the fact that the \( U(1)_A \) symmetry is broken through the QCD axial anomaly. The recent data on \( K_L \rightarrow \gamma e^+ e^- \), when combined with the branching ratio of \( K_L \rightarrow \gamma \gamma \), provide us with important information on \( \zeta \equiv a(K_2 \eta)/a(K_2 \pi^0) \), \( \zeta' \equiv a(K_2 \eta')/a(K_2 \pi^0) \) and \( \delta_p \), or equivalently, on \( \delta_n \) and \( \delta_p \). We will have two solutions for \( \delta_n \), each of which corresponds to \( \pi^- \) and \( \eta^- \) dominance in \( K_L \rightarrow \gamma \gamma \), respectively. Then, \( \delta_p \) is constrained to some region for each \( \delta_n \). In \( K^+ \rightarrow \pi^+ e^+ e^- \), we could determine \( C_7 = -0.01^{+0.03}_{-0.04} \) or \(-0.61^{+0.03}_{-0.04} \) from the best fit to \( B(K^+ \rightarrow \pi^+ e^+ e^-) \). This parameter measures the short distance contribution of the electromagnetic penguin diagram to the process \( K \rightarrow \pi l^+ l^- \). The twofold ambiguity can be lifted by the spectrum measurement of \( K^+ \rightarrow \pi^+ e^+ e^- \) at low \( m_{ee} \). For either value of \( C_7 \), we are led to rather large branching ratios for the indirect CP-violating \( K_L \rightarrow \pi^0 e^+ e^- \) at the level of a few parts in \( 10^{-10} \).

### 4.2 \( b \rightarrow s \gamma \) in MSSM and flipped SU(5) SUGRA\(^{30} \)

The purpose of this subsection is to study \( \text{BR}(b \rightarrow s \gamma) \) in two supergravity models: (i) the minimal \( SU(5) \) supergravity model\(^{31} \), and (ii) the no-scale flipped \( SU(5) \) supergravity model\(^{32} \). We find that the new CLEO bound\(^{33} \), i.e. \( \text{BR}(b \rightarrow s \gamma) < 5.4 \times 10^{-4} \) at 95\% CL, does not yet constrain the minimal \( SU(5) \) model, while some as-yet-mild constraints are imposed on the flipped \( SU(5) \) model, where a new phenomenon can drastically suppress the \( b \rightarrow s \gamma \) amplitude. We present the results for \( \text{BR}(b \rightarrow s \gamma) \) in these two models and show that improved sensitivity could probe them in ways not possible at present collider experiments.

We use the following expression\(^{34} \) for the branching ratio \( b \rightarrow s \gamma \)

\[
\frac{\text{BR}(b \rightarrow s \gamma)}{\text{BR}(b \rightarrow e e \nu)} = \frac{6\alpha}{\pi} \frac{[\eta^{16/23} A_{\gamma} + \frac{8}{3}(\eta^{14/23} - \eta^{16/23}) A_{\gamma} + C]}{I(m_c/m_b) \left[ 1 - \frac{2}{3\pi} \alpha_s(m_b) f(m_c/m_b) \right]^2},
\]

(37)

where \( \eta = \alpha_s(M_Z)/\alpha_s(m_b) \), \( I \) is the phase-space factor \( I(x) = 1 - 8x^2 + 8x^6 - x^8 - 24x^4 \ln x \), and \( f(m_c/m_b) = 2.41 \) the QCD correction factor for the semileptonic decay. The
$A_\gamma, A_g$ are the coefficients of the effective $bs\gamma$ and $bsg$ penguin operators evaluated at the scale $M_Z$. Their simplified expressions are given in the Appendix of Ref. 34, where the gluino and neutralino contributions have been justifiably neglected$^{35}$ and the squarks are considered degenerate in mass, except for the $\tilde{t}_{1,2}$ which are significantly split by $m_t$. This is a fairly good approximation to the actual result obtained in the two supergravity models we consider here, since $m_{\tilde{q}} > 200$ GeV in these models. To include at least partially the QED corrections to the $b \to s\gamma$ operator, in Eq. (37) we take $\alpha = \alpha(m_b) = 1/131.2$. We use the 3-loop expressions for $\alpha_s$ and choose $\Lambda_{\text{QCD}}$ to obtain $\alpha_s(M_Z)$ consistent with the recent measurements at LEP. For the numerical evaluation of Eq. (37) we take $\alpha_s(M_Z) = 0.118$, $\text{BR}(b \to c\ell \bar{\nu}) = 10.7\%$, $m_b = 4.8$ GeV and $m_c/m_b = 0.3$.

The subject of the QCD corrections to $\text{BR}(b \to s\gamma)$, i.e., the origin of the $\eta$ factors and the $C$-coefficient in Eq. (37), has received a great deal of attention over the years in the SM. We use the leading-order QCD corrections to the $b \to s\gamma$ amplitude when evaluated at the $\mu = m_b$ scale, i.e., $C = \sum_{i=1}^{8} b_i \eta^d_i = -0.1766$ for $\eta = 0.548$, with the $b_i, d_i$ coefficients given in Ref. 34. The result follows from the renormalization-group scaling from the scale $M_Z$ down to $\mu = m_b$ of the effective $b \to s\gamma$ operators at $M_Z$, which include the usual electromagnetic penguin operator plus some four-quark operators. Scaling introduces operator mixing effects (as exemplified by the appearance of the gluonic penguin operator in Eq. (37)) and the scaled coefficients at $\mu = m_b$ are given by linear combinations of the coefficients at scale $M_Z$. It has been pointed out$^{36}$ that the low-energy mass scale $\mu$ affects the leading-order results significantly. The combined effect on $\text{BR}(b \to s\gamma)$ of the uncertainties in $\mu$ and $\Lambda_{\text{QCD}}$ has been estimated$^{34}$ to be $\lesssim 25\%$. Recently a partial next-to-leading order calculation of the QCD effects has appeared$^{37}$, which gives somewhat smaller enhancement factors than the complete leading-order result. A full next-to-leading order QCD calculation should decrease the above mentioned uncertainties significantly.

The two models we consider here are built within the framework of supergravity with universal soft-supersymmetry breaking. The renormalization-group scaling from the unification scale down to low energies plus the requirement of radiative electroweak symmetry breaking using the one-loop effective potential, reduces the number of parameters needed to describe these models down to just five: $m_t$, $\tan \beta$, and three soft-supersymmetry breaking parameters ($m_{1/2}, m_0, A$). The sign of the superpotential Higgs mixing term $\mu$ remains as a discrete variable. The models we consider belong to this general class of models but are further constrained making them quite predictive.

The minimal SU(5) supergravity model$^{31}$ is strongly constrained by the proton lifetime and the cosmological constraint of a not too young universe$^{38-40}$. A thorough exploration of the parameter space, including two-loop gauge coupling unification yields a restricted region of parameter space to be subjected to further phenomenological tests. We find

$$2.3 \times 10^{-4} < \text{BR}(b \to s\gamma)_{\text{minimal}} < 3.6 \times 10^{-4},$$

(38)
for $\mu > 0 (\mu < 0)$, which are all within the new CLEO bound. We also obtain

$$0.90 \ (0.97) < \frac{\text{BR}(b \rightarrow s\gamma)_{\text{minimal}}}{\text{BR}(b \rightarrow s\gamma)_{\text{SM}}} < 1.20 \ (1.13), \quad (39)$$

for $\mu > 0 (\mu < 0)$. This implies that $\text{BR}(b \rightarrow s\gamma)$ would need to be measured with better than 20% accuracy to start disentangling the minimal $SU(5)$ supergravity model from the SM. Moreover, for $\mu < 0$ there is a band of points which will be difficult to tell apart (requires < 1% accuracy). It is interesting to remark that an analogous set of plots versus $m_t$ instead, does not reveal any particular structure in $m_t$. In other words, in the minimal $SU(5)$ supergravity model, a measurement of $\text{BR}(b \rightarrow s\gamma)$ within the range of Eq. (39), is unlikely to shed much light on the value of $m_t$. Of course, if the measurement falls outside this range, the model would be excluded.

The no-scale flipped $SU(5)$ supergravity model\textsuperscript{32}) assumes that the parameters $m_0$ and $A$ vanish, as is typically the case in no-scale supergravity models\textsuperscript{41}). This constraint reduces the dimension of the parameter space down to three. The choice of gauge group (flipped $SU(5)$) is made to make contact with string-inspired models which unify at the scale $M_U \sim 10^{18}$ GeV. We also consider a “strict no-scale” scenario where in addition the universal bilinear soft-supersymmetry breaking scalar mass parameter $B$ vanishes. This constraint reduces the dimension of the parameter space down to two, since $\tan \beta$ can now be computed as a function of $m_{1/2}$ and $m_t$. An interesting consequence of this scenario is that for $\mu > 0$ only $m_t \lesssim 135$ GeV is allowed, whereas for $\mu < 0$ only $m_t \gtrsim 140$ GeV is allowed. Moreover, for $\mu > 0$ the calculated value of $\tan \beta$ can be double-valued. In what follows we take $m_t = 100, 130, 140, 150, 160$ GeV for $\mu > 0 (\mu < 0)$. The results are as follows: (i) if $m_t \gtrsim 140$ GeV then $\mu < 0$ and $\text{BR}(b \rightarrow s\gamma) \gtrsim 3.5 \times 10^{-4}$; and (ii) if $m_t \lesssim 135$ GeV then $\mu > 0$ and a wide range of $\text{BR}(b \rightarrow s\gamma)$ values are possible. In both cases, with sufficiently accurate measurements, one should be able to pin down the value of the chargino mass (with a possible two-fold ambiguity) and therefore the whole spectrum of the model.

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