On A New Class of Models for Soft CP Violation

David Bowser-Chao\textsuperscript{(1)}, Darwin Chang\textsuperscript{(2,3)}, and Wai-Yee Keung\textsuperscript{(1)}

\textsuperscript{(1)}\textit{Physics Department, University of Illinois at Chicago, IL 60607-7059, USA}
\textsuperscript{(2)}\textit{NCST and Physics Department, National Tsing-Hua University, Hsinchu 30043, Taiwan, R.O.C.}
\textsuperscript{(3)}\textit{Institute of Physics, Academia Sinica, Taipei, R.O.C.}

Introduction

Recently, two new classes of models of soft CP violation have been proposed as alternatives to the standard KM model. The models aim at reproducing the attractive characters of KM model and at the same time solve the long-standing strong CP problem.

These models share the features of imposing soft (or spontaneous) CP violation in order to avoid tree level KM phase as well as the tree level strong CP phase. To account for the measured value of the CP-violating parameter $\epsilon$ of the neutral kaon system, the models employ a new heavy sector of scalars and vectorial fermions. The first class of models, that we shall roughly classified as the right-handed (RH) models, uses a heavy sector that couples only to right-handed down type quarks, which are $SU_L(2)$ singlets. Instead, the alternative left-handed ( LH) models uses new particles that couple only to the ordinary left-handed quarks, which are $SU_L(2)$ doublet.

In this report, we review the phenomenology of the RH models and make comments on the LH models as we go for comparison. The more extensive presentation is under preparation. Note that the ideas in these directions were presented by Barr and his collaborators some time ago.

The heavy sector of a typical model for our purposes requires two additional Higgs singlets, $h_\alpha (\alpha = 1, 2)$ of charge $q_h$ and a vectorial pair of heavy fermions, $Q_{L,R}$, of electromagnetic charge $-\frac{1}{2} + q_h$. One can also choose to have two pairs of fermions and only one heavy Higgs singlet. In addition, one can choose to assign the scalars or the fermions to be carrying the color such that together they will couple to the right-handed down type quarks, $d_{Ri}$. In case of neutral, colorless fermions (the neutrinos), it is not even necessary to have vectorial pairs. One can also use the freedom in choosing charge $q_h$ to avoid fractionally charged hadrons. Most of the phenomenology mention below are more or less independent of the choice of $q_h$ and color. Relevant new terms in the Lagrangian are:

\begin{equation}
\mathcal{L}_{h_1} = [(g\lambda_{\alpha}(Q_{L}d_{R})h_\alpha + M_Q Q_{L}Q_{R}) + h.c.]
- (m^2)_{\alpha\beta}h_\alpha^\dagger h_\beta
- \kappa_{\alpha\beta}(\phi^\dagger \phi - |\langle \phi \rangle|^2) h_\alpha^\dagger h_\beta
- \kappa_{\alpha\beta\gamma\delta}h_\alpha^\dagger h_\beta h_\gamma^\dagger h_\delta,
\end{equation}

where $\phi$ is the Standard Model Higgs doublet. The soft breaking of CP symmetry implies a special basis where $\lambda, c, k'$ and the SM Yukawa couplings are real. If fermions carry color, we also require (see below) that dim-3 couplings, $M_Q$, are real to avoid tree level contribution to $\theta$. This leaves only a single CP violating parameter: $\text{Im}(m^2)_{12}$. We can diagonalize $(m^2)_{\alpha\beta}$ by a unitary matrix $U_{ai}$ which in general is complex: $h_\alpha = U_{ai}H_i$, with $H_i$ the mass eigenstates. The quark-Higgs interaction in the mass eigenstate basis becomes

\begin{equation}
\mathcal{L}_{Q_{HI}} = g \sum_{q=d,s,b} \xi_{qj} \bar{Q}_{L}q_{R} H_j + \text{h.c.,}
\end{equation}

where $H_i = U_{ai}h_\alpha$ is the mass eigenstates, $\xi_{qj} \equiv \lambda_{qj} U_{aj}$. The Yukawa couplings $\xi_{qj}$ are defined relative to the gauge coupling $g$ of $SU_L(2)$. The rephasing-invariant measure of CP violation are

\begin{equation}
A_{ij}^{qq'} = \lambda_{qj} \lambda_{q'j}^* U_{bij} U_{aij} = \xi_{qj}^* \xi_{q'j},
\end{equation}

with $(q, q' = d, s, b)$ and $i, j = 1, 2$. For flavor conserving amplitudes like EDM, we define the counterpart $B_{ij} = \kappa_{ij} U_{bij}^* U_{aij}$. $A_{ij}^{qq'}$ as well as $B_{ij}$ are Hermitian in indices $i, j$. Thus CP is broken only in the off-diagonal terms. As a result, they contribute to CP violation only when both the light and heavy charged Higgs are involved in a diagram at the lowest order. It is also possible to avoid the CP violating part of the coupling $\kappa$ if one chooses to have only one heavy Higgs singlet and...
breaks CP symmetry in the dimension-3 heavy fermion mixing terms instead as long as these heavy fermions are colorless. The contribution related to parameter $\kappa'$ in the last term of (1) occurs only at the higher loop level and generally can be ignored.

Before continuing, we like to emphasize again that if the new fermions carry color it is necessary to impose CP symmetry on their bare masses also in order to avoid tree level strong CP problem. When there are more than one pairs of vectorial fermions, one can potentially make the mass matrix Hermitian, however that would require some additional symmetry to be implemented.

**Constraint from $\epsilon$**

With CP broken only softly, the CKM matrix is real at tree level. Leading contribution to the CP violating parameter $\epsilon$ is due to the box diagram with only heavy particles in loop.

$$H^{\Delta S=2} = \frac{G_F^2 m_W^2}{16\pi^2} \sum_{i=R,L} C_i^{\Delta S=2}(\mu) O_i^{\Delta S=2}(\mu), \quad (4)$$

with $O_i^{R,L=2} = \bar{s}_\mu(1 + \gamma_5)d \bar{s}_\mu(1 + \gamma_5)d$. \quad (5)

The $W^\pm$ diagrams yield a purely real Wilson coefficient $C_i^{\Delta S=2}(\mu)$; CP violation is due solely to the operator $O_i^{R,L=2}$ rather than $O_i^{L,R=2}$, in contrast to the KM model, because the complex coefficient $C_i^{R,L=2}(\mu)$ is generated by the charged Higgs. At the scale $\mu = M_Q$, we have

$$C_i^{\Delta S=2}(M_Q) = 2 \xi_d \xi_s^* \xi_d \xi_s^* \frac{m_W^2}{M_Q^2} \frac{f(x_2) - f(x_1)}{x_2 - x_1}$$

$$+ \sum_{i=1,2} (\xi_d \xi_s^*)^2 \frac{m_W^2}{M_Q^2} \frac{df(x_i)}{dx}, \quad (6)$$

with $f(h) = (1 - h)^{-2}(1 + 2h + h^2 + h^2 \ln h)$.

The real part of the diagram can contribute to part of the $K_L - K_s$ mass difference while the imaginary gives rise to $\epsilon$. This is analyzed in detail in Ref.\cite{3}. For illustration here, we can take the (decoupling) limit $m_2 \gg m_1$ and assume $m_1 = M_Q$ for simplicity. Demanding that the imaginary part of the box diagram gives enough contribution to $\epsilon$ while the corresponding real part gives just a fraction $F$ of the mass difference $\Delta m_K$, we obtain the constraints,

$$\text{Im} \left( A_{sd}/(0.058)^2 \right) R_{Q}^2 = 1, \quad (7)$$

$$\text{Re} \left( A_{sd}/(0.058)^2 \right) R_{Q}^2 = 156F. \quad (8)$$

where $R_Q = 300 \text{ GeV}/M_Q$. The reasonable constraint $|F| < 1$ can be easily satisfied. It is important to emphasize that, in RH models, the heavy particle box diagrams induce a right-handed four fermion operator, contrary to the LH models and the KM model in which the leading CP violating operators are left-handed.

**Constraints from $(\epsilon'/\epsilon)$**

The leading contribution to the direct CP violating parameter $\epsilon'$ is due to the gluonic penguin diagrams with only the heavy particles in the loop.

At the electroweak scale, the Wilson coefficient is

$$\tilde{C} = -\alpha_s \sum_i \xi_d \xi_s^* \frac{m_W^2}{M_Q^2} F \left( \frac{m_{H_u}}{M_Q^2} \right),$$

$$F(h) = \left[ \frac{(2h - 3)h^2 \ln h}{(1 - h)^4} + \frac{7 - 29h + 16h^2}{6(1 - h)^3} \right].$$

The electromagnetic penguin and long distance effects can contribute to the imaginary part of the $\Delta I = 3/2$ amplitude and give a small contribution. This is analyzed in detail in Ref.\cite{3}. In our decoupling limit, the result is

$$\epsilon'/\epsilon = -1.9 \times 10^{-5} \text{Im}(A_{sd}/(0.058)^2) R_{Q}^2. \quad (9)$$

$$= \pm 1.9 \times 10^{-5} \left( \sqrt{(156F)^2 + 1 - 156F} \right) \frac{R_{Q}}{\sqrt{F}},$$

using the constraints in Eq.\cite{3}. For $R_Q = 1$ and $F \approx 0$ (or $-0.3$), $\epsilon'/\epsilon = 1.4 \times 10^{-5}$ (or $1.3 \times 10^{-4}$), which is roughly the same order of magnitude as the KM model. One can of course makes the model more superweak\cite{1} by setting the scale higher (and $R_Q$ smaller) and $A_{sd}$ larger. For $A_{sd}$ of order one, $M_Q$ is roughly 100 TeV.
Constraints from Strong CP $\theta_{QCD}$ and the induced KM phase

In both RH and LH models, $\theta_{QCD}$ is only induced starting at the two-loop level, via generation of complex down-flavor quark masses as long as the coupling $\kappa$ exists. A typical diagram is shown in Fig. 1 in Ref. 1, this effect does not require more than one flavor of down-quark. However, it does require both charged Higgs to be involved.

\[
\begin{align*}
\langle \phi \rangle & \quad \xrightarrow{H_1} \quad \xrightarrow{H_2} \quad \phi \\
\imath \, m_1 & \quad \xrightarrow{d_R} \quad Q_L \quad \xrightarrow{d_R} \quad d_L
\end{align*}
\]

Roughly, $\theta_{QCD} \sim g^2 I \, \text{Im}(A_{12}^d B_{12})/(16\pi^2)^2$. The factor $I$, of order one, is given by the integral

\[
I = \frac{1}{\theta} \int_0^1 \frac{dz}{1-z} \int_0^{1-z} \frac{dx}{(1-z)M_{\phi^0}} \left[ \begin{array}{c}
\log \frac{z(1-z)M_{\phi^0}}{z^2M_{\phi^0} + xM_{\phi^0} + yM_{\phi^0}^2} - (m_{H_1} \leftrightarrow m_{H_2}) \\
\end{array} \right],
\]

with the Feynman parameters $y = 1-x-z$. The integral $I$ vanishes at the degenerate case $m_{H_2} = m_{H_1}$, but its size approaches 1 as $m_{H_2} \to \infty$. This non-decoupling phenomena is not surprising because, in the large $m_{H_2}$ limit, the CP is a broken symmetry. Numerically, the present constraint, $\theta_{QCD} < 10^{-8}$, can easily be accommodated. In addition, there are also three loop diagrams due to the gluonic contribution. A typical graph is shown below.

This contribution is independent of the coupling $\kappa$. However, unless $\kappa$ happens to be very small, it may not be competitive with the two loop contributions because of the KM angle suppression and the additional loop factor. Both the 2-loop and the 3-loop contributions also exist generically in LH models. However, they are typically numerically smaller in that case. The two loop contributions disappear (in both RH and LH models) of course when only one complex scalar boson is used as mentioned earlier.

Since CP is broken at higher energy, a non-vanishing KM phase $\eta$ (defined in the Wolfenstein parametrization) can in general be loop induced. It can originate either from the loop-induced complex mass matrix or from that of the complex kinetic energy terms. The contribution from complex mass matrix is similar to the analysis of $\theta$ and is therefore small. The contribution from complex kinetic terms is induced at the one loop level in LH models as analyzed in Ref. 1. It is in general also suppressed by some small KM mixing angles and small mass ratios and therefore numerically tiny enough to be ignored phenomenologically. In contrast, in RH models, since the KM phase is related only to the rotation of the left-handed quarks, such contribution will not arise until at the two loop level as given in the following figure. Therefore induced $\eta$ is even smaller.

Constraints from electric dipole moments

A down-flavor quark EDM, however, is generated at the two-loop (or three-loop) level, in parallel with the generation of complex down quark masses discussed above. The typical contribution is given by diagrams similar to those for $\theta$, except with an external photon attached to internal charged lines. An estimate of the two loop contribution gives EDM which is consistent with the current experimental bound. The contribution from chromo-electric dipole moment of gluon (the Weinberg operator) won’t arise until three loop level (even with $\kappa$ coupling) and therefore expected to be small. The electron couples only indirectly with the CP violating sector, so its EDM vanishes at two loops and the three loop contributions are insignificantly small.

$B^0\bar{B}^0$ Mixing, $b \to s\gamma$ and Other Constraints

Another (much weaker) constraint to be considered is that from the $B^0\bar{B}^0$ mass splitting. Using the usual estimate of strong form factors involved and the experimental value for $\Delta M_{B^0}$, we have

\[
\delta(\Delta M_{B^0})/\Delta M_{B^0} = 1.1 \times 10^{-3} R_{Q}^2 \text{Re} \left( A_{bd}/0.058^2 \right)^2,
\]

so even taking $A_{bd} = (0.15)^2$, the fractional contribution is only about 5% for $M_Q = 300$ GeV.
In RH models, the operator induced by the exotic sector has helicity opposite to the SM contribution, the two do not interfere in the rate. In the decoupling limit with \( F_{b \to s \gamma} \equiv \delta B(b \to s \gamma)/B(b \to s \gamma)_{\text{SM}} \), we have

\[
F_{b \to s \gamma} = 6.4 \times 10^{-6} \left| \frac{R^2_{Q}}{V_{bb} V_{ts}^*} \cdot \frac{A_{bs}}{(0.058)^2} \right|^2.
\]

Furthermore, the relevant parameter \( A_{bd} \) is not subjected to constraints from \( \epsilon \) or \( \epsilon' \). If it is of the same size as \( A_{sd} \), the deviation from the SM would be negligible and the future B factory would observe only a collapsed unitarity KM angle. However, if \( A_{bd} \gg A_{sd} \), the triangle can looks substantially different from that predicted by standard KM model. It is worthwhile to point out that, in LH models, the exotic contribution gives rise to operator that will interfere with that of SM and therefore the model is more severely constrained.

### Decays of New Particles

In the generic RH model, \( h \) and \( Q \) can be assigned a new conserved quantum number which guarantees a stable lightest exotic particle, either \( H_1 \) or \( Q \). However, when \( q_b = -1 \), an interaction, \( h_a L_i L_j \), is allowed, which can lead to \( H^- \) (on-shell or off-shell) decays into \( l^- \nu \). Even so, lepton number is still conserved, just as in the SM, since \( Q \) and \( H \) will naturally carry lepton number (\( L = \pm 2 \)). Another way for \( H \) to decay is to introduce a second Higgs doublet and let \( H \) couple to two different Higgs doublets. In that case \( H \) can decay into a neutral Higgs, plus a charged Higgs which in turn decays into ordinary quarks and leptons.

### Spontaneous broken CP symmetry

To show how the above softly broken CP symmetry can in fact originate from a spontaneously broken one, one can first add a CP-odd scalar, \( a \), which develops a non-zero vacuum expectation value (VEV) and breaks CP. However, this scalar will in general couple to \( Q_i Q_R \) and give rise a complex tree level \( M_Q \) and, therefore, a tree level \( \theta_{QCD} \). To avoid this, one can add another scalar singlet, \( s \), which is CP-even and impose a discrete symmetry which changes the signs of both \( a \) and \( s \) and nothing else. As a result, a term such as \( ia\bar{Q}\gamma_5 Q \) is forbidden.

The only additional term relevant for CP violation is \( i \left[ s a (h_1^\dagger h_2 - h_2^\dagger h_1) \right] \) which generates a complex \( (m^2)_{12} \) after both \( s \) and \( a \) develop VEVs. The extra neutral Higgs bosons will mix with the SM Higgs, but since \( s \) and \( a \) do not couple to fermions directly, they have tiny scalar-pseudoscalar coupling to fermions only at loop level. As a result, the CP phenomenology considered below applies equally well to both the softly and spontaneously broken versions of our model.

### Interplay between Strong and Weak CP Phase

The RH models provide a good example to look at the subtlety, raised in Ref. 1, involving the interplay of the CP phases between the strong and the weak interactions. For our purposes, we can focus on the reduced effective theory which contains CP-conserving Standard Model-type interactions and vanishing \( \theta_{QCD} \), plus the new, induced superweak interaction with the strength \( C_{\Delta S=2}^R \) defined in Eq. (4). We shall consider, for the sake of argument, the scenario in which the up quark is massive while \( m_d \) is zero. Without the new \( C_{\Delta S=2}^R \) interaction the parameter \( \theta_{QCD} \) is then unphysical, with CP a good symmetry. (By the usual argument, with a massless quark present — in this case, the \( d \) quark — the right-handed component of that quark can be rotated to absorb \( \theta_{QCD} \) via the axial anomaly, while otherwise leaving the lagrangian invariant.) With the addition of the induced interaction \( H^2 S^{-2} \), however, \( \theta_{QCD} \) becomes physical, as can be seen by considering the following cases:

(a) \( C_{\Delta S=2}^R \) is complex and \( \theta_{QCD} = 0 \), CP is violated. In this case, the correct (non-zero) value for \( \epsilon \) can be calculated without complication; all hadronic matrix elements (modulo absorptive contributions) can correctly be assumed to be real. It is illuminating to consider the calculation of \( \epsilon \) in another basis, which is obtained by a phase rotation of \( d_R \) such that \( C_{\Delta S=2}^R \) becomes real and \( \theta_{QCD} \) non-zero.

Since the two theories are the same, one must arrive at the same result for \( \epsilon \). One can thus draw a rather surprising conclusion: \( \theta_{QCD} \) can also, in certain situations, contribute to \( \epsilon \).

In fact, from the way we obtain the \( \theta_{QCD} \) contribution to \( \epsilon \) in this example, one realizes that there is an important subtlety here. The actual contribution from \( \theta_{QCD} \) to \( \epsilon \) is correlated to the explicit mechanism of CP violation, which in our current example is the superweak \( C_{\Delta S=2}^R \). A related result is that when \( \theta_{QCD} \) is not zero, how each hadronic matrix element develops a phase also depends on the particular electroweak mechanism of CP violation in the theory. In the present case, the CP violating coupling also happens to be the chiral symmetry breaking phase.

Another lesson one learns is that the usual argument which concludes that the contribution of \( \theta_{QCD} \) to CP-violating quantities such as the neutron electric dipole moment (edm) must be proportional to \( m_u m_d \), is not strictly correct, a counterexample is offered by the simplified model presented above. The role of \( m_d \) is replaced by the coupling \( C_{\Delta S=2}^R \). Of course, \( C_{\Delta S=2}^R \) breaks the chiral symmetry associated with \( d \) quark, so that the \( d \) quark will certainly pick up mass at some (probably...
higher-loop) level, but the point is that the \( C^R_{\Delta S=2} \) coupling plays a much more direct role in the contribution of \( \theta_{QCD} \) than even the induced \( m_d \)!

Now we come to an apparent paradox whose resolution gives even further insight into the interplay between strong and weak CP phases.

We parenthetically noted above that if redefinition of the quark phases generates an imaginary part for the quark mass matrix not proportional to the identity matrix, the low-energy meson states must be suitably reinterpreted to ensure stability of the vacuum around which we carry out perturbation theory. This redefinition explicitly reintroduces the phase(s) rotated from the couplings into certain hadronic matrix elements, to ensure rephasing invariance. If both \( m_d, m_u \) vanish, then arbitrary rotation of the corresponding right-handed quarks seems to have no effect on vacuum stability, since the mass matrix is left real and diagonal (only \( m_u \) non-zero). Then, apparently, all phases may be arbitrarily rotated away, and with them, any possibility of CP violation. Specifically, consider the following variant of the two cases already considered:

(b) Let \( C^R_{\Delta S=2} \) be complex, but take \( \theta_{QCD} \) to be zero. If both \( m_u, m_d \) are strictly zero, is CP conserved or violated? At first glance, one might claim the phase in \( C^R_{\Delta S=2} \) to be unphysical, since a combined phase rotation of the form \( u_R \rightarrow e^{-i\theta} u_R \) and \( d_R \rightarrow e^{i\phi} d_R \) can make \( C^R_{\Delta S=2} \) real and maintain \( \theta_{QCD} = 0 \). It is very tempting to claim that CP violation is proportional to \( m_u \) for a small up quark mass and further that there is no CP violation when \( m_u \rightarrow 0 \) because the phase of \( C_{\Delta S=2} \) then becomes absorbed.

This conclusion is incorrect, however, because we have ignored the vacuum degeneracy in the case of massless \( u \) and \( d \) quarks. Different choices of vacua would give different CP violation. It is true that there exists one very special vacuum where CP is conserved. However, a general vacuum possesses chiral condensate with a phase uncorrelated to that of \( C^R_{\Delta S=2} \), and thus CP violation usually occurs, even if \( C^R_{\Delta S=2} \) is real (since what is important is the relative phase between the vacuum and \( C^R_{\Delta S=2} \)). This idea can be demonstrated directly in the chiral effective lagrangian approach. The chiral field \( \Sigma(3 \times 3 \text{ unitary matrix}) \) can be perturbed around a vacuum configuration \( \text{diag}(e^{-i\phi}, e^{i\phi}, 1) \). If \( C^R_{\Delta S=2} \) is turned off, the strong interaction is independent of \( \phi \) because of the chiral symmetry. However, with \( C^R_{\Delta S=2} \), the phase \( \phi \) has physical meaning and has implications with respect to CP violation. Now we include effects of the real quark masses \( m_u \neq 0 \) and \( m_d \neq 0 \). Their net effect is simply to pick out a particular vacuum, with \( \Sigma = \text{diag}(1,1,1) \). In this case of vacuum alignment, a complex \( C^R_{\Delta S=2} \) is necessary, but also sufficient, for CP violation, since again it is the relative phase between \( C^R_{\Delta S=2} \) and the vacuum that is important. In some sense, the (possibly infinitesimal) up and down quark masses enforce CP violation, in the particular case that \( C^R_{\Delta S=2} \) is real, whereas in the massless case, CP violation is still generally expected via the vacuum phase.

**Conclusion**

We have reviewed a new (RH) class of models whose CP violation is solely mediated by exotic Higgs bosons and fermions that couple to the right handed quarks, and compared with another (LI) class of models whose exotic particles couple to the left-handed quarks. The model naturally prevents the strong CP problem. Both classes of models are surprisingly similar to the KM model in the sense that the CP-breaking mechanism is seemingly milliweak (if the exotic particle scale is chosen to be as low as possible), while its phenomenology (as studied here) is quite superweak-like. The phenomenological distinction between the two will likely be made clear in experiments planned for the B factory. A careful and detailed analysis of such issues is clearly necessary and is in progress.

D. B.-C. and W.-Y. K. are supported by a grant from the DOE of USA, and D. C. by the NSC of R.O.C. We thank H. Georgi, S. Glashow, P. Frampton, R. Mohapatra, and L. Wolfenstein for very useful discussions.

**References**

1. D. Bowser-Chao, D. Chang, and W.-Y. Keung, Phys. Rev. Lett. 81, 2028 (1998).
2. H. Georgi and S. Glashow, hep-ph/9807399. See also, P. Frampton and M. Harada, hep-ph/9809402.
3. M. Kobayashi and T. Maskawa, Prog. Theor. Phys. 49, 652 (1973).
4. D. Bowser-Chao, D. Chang, and W.-Y. Keung, in preparation.
5. S. Barr and A. Zee, Phys. Rev. Lett. 55, 2253 (1985); J. Nieves, Nucl. Phys. B189, 189 (1981); Phys. Lett 164B, 85 (1985); S. Barr, Phys. Rev. D34, 1567 (1986); S. Barr and E. M. Freire Phys. Rev. D41, 2129 (1990).
6. L. Wolfenstein, Phys. Rev. Lett. 13, 562 (1964).
7. L. Wolfenstein, Phys. Rev. Lett. 51, 1945 (1983); L.–L. Chau and W.–Y. Keung, Phys. Rev. Lett. 53, 1802 (1984).
8. S. Mele, hep-ph/9810333; see also R. Barbieri, L. Hall, A. Stocchi, and N. Weiner, Phys. Lett. B425 119 (1998) or hep-ph/9712252.
9. D. Bowser-Chao, D. Chang, and W.-Y. Keung, Chin. J. Phys. 35, 842 (1997), hep-ph/9803275.