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SUPER KOBAYASHI-MASKAWA CP-VIOLATION

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ABSTRACT

We argue that in the minimal SUSY extension of the standard model, CP-violation cannot be explained through SUSY phases alone. But SUSY graphs, especially with gluinos, can make important contributions to CP-violation through the Kobayashi-Maskawa phase.

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The SU(2) x U(1) gauge symmetry of the standard model[1] for electroweak interactions is spontaneously broken by the (real) vacuum expectation value of a single scalar doublet. This implies that in this minimal framework, CP-violation only arises from complex Yukawa couplings while flavour changing neutral processes are naturally forbidden at the tree level. As a consequence, CP violating flavour changing neutral processes are induced through charged weak gauge boson exchanges and described by the Kobayashi-Maskawa (KM)[2] mixing matrix. In particular, assuming a not-too-light top quark, the estimated ε parameter associated with the $K^0 - \bar{K}^0$ system is compatible with the experimental value.[3] Flavour changing neutral processes in general and the $K^0 - \bar{K}^0$ system in particular are, in fact, stringent tests for any extension of the presently successful standard model.

Recently, a minimal supersymmetric version of this model[4] based on N=1 supergravity has been proposed. Although two scalar doublets are now needed, flavour changing neutral currents are still forbidden at the tree level, due to supersymmetry (SUSY). However, complex soft breaking terms such as gaugino Majorana masses or trilinear Yukawa-like scalar self-interactions [5], low energy remnants of the super Higgs mechanism [6] responsible for the breaking of local supersymmetry, introduce new physical phases in the game. Moreover, with a typical breaking scale of the order of 10^{10} Gev, large renormalization effects are expected. In particular ,it has been stressed that the new gauge couplings between squark ($\tilde{q}$), quark (q) and gaugino ($\tilde{g}$)obtain an important flavour changing piece from one loop radiative corrections induced by the charged scalars. The gluino, in particular, could produce a large contribution to ε because it is associated with the strong coupling constant. We therefore focus our attention on these contributions. However,using restrictions from the experimental bound on the electric dipole moment of the neutron (EDMN)[7] we find that SUSY-phases alone cannot be responsible for CP-violation.

On the other hand, we show that CP-violation from the KM-phase through supersymmetric diagrams can be of the right order of magnitude. This becomes especially important if the top quark mass turns out to be relatively small. Indeed, in this case , conventional non-supersymmetric CP-violation might not be sufficient[3]. Furthermore, SUSY-graphs bring the KM-contributions to the EDMN nearer to the experimental bound. Thus even if the intrinsic SUSY-phases are small, SUSY can play an important rôle in CP-violation.

From this point of view, we derive restrictions on the squark and gluino masses coming from the ε parameter. As a byproduct, from the real part of the $K^0-\bar{K}^0$ matrix element, we only get mild constraints.

In summary, the minimal SUSY extension of the standard model successfully emerges from this preliminary but usually dangerous entrance examination.

In this model [4),5)], the mass matrix for the down squarks ($\tilde{d}$) is related to the quark mass matrices M_d, M_u in the following way:

$$M^2_{\tilde{d}} \simeq \begin{pmatrix} \mu_L^2 1 + M_d M_d^\dagger + c M_u M_u^\dagger & A^* m_{3/2} M_d \\ A m_{3/2} M_d^\dagger & \mu_R^2 1 + M_d^\dagger M_d \end{pmatrix} \tag{1}$$

where A is the typical complex soft breaking parameter induced by the "hidden" superpotential responsible for the super Higgs mechanism

$$A = |A| e^{-2i\varphi_A} \tag{2}$$

while the $\mu_{L,R}$ mass parameters are of the same order as the gravitino mass $m_{3/2}$. The coefficient c, which is of order 1[8),9)], will be crucial in our discussion. It is associated with the main one loop flavour violating correction to $M^2_{\tilde{d}}$ and thus implies a CP-violating flavour changing piece in the new gauge interactions for down quarks. For the gluino ($\tilde{g}$)[10)] this interaction is given by

$$\mathcal{L}_{\tilde{g}\tilde{d}d} = i\sqrt{2}\, g_s \tilde{d}_i^\dagger \bar{\tilde{g}}_a T^a \left\{ \Gamma_L P_L + \Gamma_R P_R \right\}^{ij} d_j$$

with

$$P_L = \frac{1-\gamma_5}{2}, \quad P_R = \frac{1+\gamma_5}{2} \tag{3}$$

where g_s denotes the strong coupling constant and T^a are the representation matrices of $SU(3)_c$. Similar Lagrangians describe the interactions of the other gauginos[10)]. The 6 x 3 matrices $\Gamma_{L,R}$ are defined as

$$\Gamma_L = e^{i\varphi_g}\, \tilde{U}^\dagger \begin{pmatrix} U_L^d \\ 0 \end{pmatrix}, \quad \Gamma_R = e^{-i\varphi_g}\, \tilde{U}^\dagger \begin{pmatrix} 0 \\ U_R^d \end{pmatrix} \tag{4}$$

where $U^d_{L,R}$ and $\tilde{U}$ diagonalize the down quark mass matrix and the down squark mass matrix respectively:

$$U_L^{d\dagger} M_d U_R^d = \hat{M}_d, \quad \tilde{U}^\dagger M^2_{\tilde{d}} U = \hat{M}^2_{\tilde{d}} = \mathrm{diag}\,(\tilde{M}_j^2) \tag{5}$$

The SUSY-phase ϕ_g appearing in the $\Gamma_{L,R}$ matrices arises from a suitable phase rotation of the gluino field whose soft breaking Majorana mass term is in general complex:

$$m_{\tilde{g}} = |m_{\tilde{g}}|\, e^{-2i\varphi_g} \tag{6}$$

For the diagonalization of the down squark mass matrix (1), we will proceed in two steps by introducing two unitary matrices $\tilde{U}_1$ and $\tilde{U}_2$ such that

$$\tilde{U} = \tilde{U}_1 \tilde{U}_2 \,, \quad \text{with} \quad \tilde{U}_1 = \begin{pmatrix} e^{i\varphi_A} U_L^d & 0 \\ 0 & e^{-i\varphi_A} U_R^d \end{pmatrix}. \tag{7}$$

We get

$$\tilde{U}^\dagger M_{\tilde{d}}^2 \tilde{U} = \tilde{U}_2^\dagger \begin{pmatrix} \mu_L^2 1 + \hat{M}_d^2 + c K^\dagger \hat{M}_u^2 K & |A| m_{3/2} \hat{M}_d \\ |A| m_{3/2} \hat{M}_d & \mu_R^2 1 + \hat{M}_d^2 \end{pmatrix} \tilde{U}_2 \tag{8}$$

where K is the KM matrix. In particular, we see that only the combination $\phi_g - \phi_A$ is relevant for CP-violation.

Let us now turn to the matrix element $M_{21}(\tilde{g}) \equiv \langle \bar{K}^0 | H_{eff}^{\Delta S=2}(\tilde{g}) | K^0 \rangle$ induced by gluino exchange (see Fig. 1). Performing a complete evaluation, we obtain:

$$\begin{aligned} M_{21}(\tilde{g}) = \frac{2\alpha_s^2}{m_{\tilde{g}}^2} \Big\{ I_{jm} \Big[& \left(\Gamma_L^{js*} \Gamma_R^{jd} \Gamma_L^{ms*} \Gamma_R^{md} + L \leftrightarrow R \right) \left(S_+^D - \tfrac{1}{2} S_+^M + \tfrac{1}{8} \Sigma^M \right) + \\ & + 2 \Gamma_L^{js*} \Gamma_L^{jd} \Gamma_R^{ms*} \Gamma_R^{md} S_-^D + \\ & + \left(\Gamma_L^{js*} \Gamma_L^{jd} \Gamma_L^{ms*} \Gamma_L^{md} + L \leftrightarrow R \right) \left(\tfrac{1}{2} V_+^M \right) \Big] + \\ + \tfrac{1}{4} K_{jm} \Big[& 2 \Gamma_L^{js*} \Gamma_R^{jd} \Gamma_R^{ms*} \Gamma_L^{md} \left(V_-^D + V_-^M \right) + \\ & + 2 \Gamma_L^{js*} \Gamma_L^{jd} \Gamma_R^{ms*} \Gamma_R^{md} \left(2 S_-^M \right) + \\ & + \left(\Gamma_L^{js*} \Gamma_L^{jd} \Gamma_L^{ms*} \Gamma_L^{md} + L \leftrightarrow R \right) V_+^D \Big] \Big\} \end{aligned} \tag{9}$$

where the terms with superscript D and M arise from diagrams of Figs. 1a and 1b respectively. The functions I_{jm} and K_{jm} are given by

$$\begin{aligned} I_{jm} &= \frac{1}{z_j - z_m} \left[\frac{z_j \ln z_j}{(1-z_j)^2} + \frac{1}{1-z_j} - (z_j \leftrightarrow z_m) \right] \\ K_{jm} &= \frac{1}{z_j - z_m} \left[\frac{z_j^2 \ln z_j}{(1-z_j)^2} + \frac{1}{1-z_j} - (z_j \leftrightarrow z_m) \right] \quad \text{with} \quad z_j = \left(\frac{\tilde{M}_j}{m_{\tilde{g}}} \right)^2. \end{aligned} \tag{10}$$

Introducing the following matrix elements

$$V_{AB}^{D(M)} = \langle \overline{K^0} | \bar{s}_\alpha \gamma^\mu P_A d_\beta \bar{s}_\gamma \gamma_\mu P_B d_\delta | K^0 \rangle (T^a T^b)_{\alpha\delta} (T^{b(a)} T^{a(b)})_{\gamma\beta}$$

$$S_{AB}^{D(M)} = \langle \overline{K^0} | \bar{s}_\alpha P_A d_\beta \bar{s}_\gamma P_B d_\delta | K^0 \rangle (T^a T^b)_{\alpha\delta} (T^{b(a)} T^{a(b)})_{\gamma\beta} \tag{11}$$

$$\Sigma_{A}^{D(M)} = \langle \overline{K^0} | \bar{s}_\alpha \sigma_{\mu\nu} P_A d_\beta \bar{s}_\gamma \sigma^{\mu\nu} d_\delta | K^0 \rangle (T^a T^b)_{\alpha\delta} (T^{b(a)} T^{a(b)})_{\gamma\beta}$$

with

$$A, B = L, R \; ; \; \alpha, \beta, \gamma, \delta = 1,2,3 \quad \text{(colour indices)}$$

then the Dirac (Majorana) terms in (9) satisfy the relations:

$$V_{+}^{D(M)} = V_{LL}^{D(M)} = V_{RR}^{D(M)} \quad ; \quad S_{+}^{D(M)} = S_{LL}^{D(M)} = S_{RR}^{D(M)}$$

$$V_{-}^{D(M)} = V_{LR}^{D(M)} = V_{RL}^{D(M)} \quad ; \quad S_{-}^{D(M)} = S_{LR}^{D(M)} = S_{RL}^{D(M)}$$

$$\Sigma^{D(M)} = \Sigma_{L}^{D(M)} = \Sigma_{R}^{D(M)} . \tag{12}$$

These matrix elements can be evaluated in the vacuum insertion method. With $\bar{K}^0(\vec{p}) = CP\, K^0(-\vec{p})$, we obtain

$$\begin{aligned}
V_{+}^{D} &= \tfrac{11}{27} N & V_{+}^{M} &= \tfrac{2}{27} N \\
V_{-}^{D} &= -(\tfrac{8}{27} + \tfrac{2}{9} R) N & V_{-}^{M} &= (\tfrac{1}{27} - \tfrac{2}{9} R) N \\
S_{+}^{D} &= -\tfrac{13}{54} R N & S_{+}^{M} &= \tfrac{5}{54} R N \\
S_{-}^{D} &= (\tfrac{1}{18} + \tfrac{8}{27} R) N & S_{-}^{M} &= (\tfrac{1}{18} - \tfrac{1}{27} R) N \\
\Sigma^{D} &= \tfrac{2}{3} R N & \Sigma^{M} &= \tfrac{2}{3} R N
\end{aligned} \tag{13}$$

with

$$N \equiv -\frac{m_K f_K^2}{2} \; , \quad R \equiv \frac{m_K^2}{(m_d + m_s)^2} . \tag{14}$$

Now let us analyze the $\tilde{g}$ contribution to the ϵ_m parameter [11] due to the SUSY phase $\phi \equiv \phi_{\tilde{g}} - \phi_A$. For simplicity, we will restrict ourselves to two generations (no KM phase) in which case we can approximately calculate the $\Gamma_{L,R}$ matrices defined in (4):

$$\Gamma_L = e^{i\varphi} \tilde{U}_2^{+} \begin{pmatrix} 1 \\ 0 \end{pmatrix} \simeq e^{i\varphi} \begin{pmatrix} c\theta_c c\rho_1 & -s\theta_c c\rho_2 \\ s\theta_c c\rho_1 & c\theta_c c\rho_2 \\ s\rho_1 & 0 \\ 0 & s\rho_2 \end{pmatrix} , \; \Gamma_R = e^{-i\varphi} \tilde{U}_2^{+} \begin{pmatrix} 0 \\ 1 \end{pmatrix} \simeq e^{-i\varphi} \begin{pmatrix} -c\theta_c s\rho_1 & s\theta_c s\rho_2 \\ -s\theta_c s\rho_1 & -c\theta_c s\rho_2 \\ c\rho_1 & 0 \\ 0 & c\rho_2 \end{pmatrix} \tag{15}$$

with

$$tg(2\rho_i) = -2|A| m_{d_i} m_{3/2} / (\mu_L^2 - \mu_R^2) \tag{16}$$

and θ_c being the Cabibbo angle. Here, $c\rho_1$, $s\theta_c$ etc. stand for $\cos\rho_1$, $\sin\theta_c$ etc. Usually $\mu_L - \mu_R$ turns out to be of the order of a few GeV[12]. From the way the ϕ-phase appears in $\Gamma_{L,R}$ we immediately conclude that only the first two terms in (9) can give an imaginary contribution to $M_{21}(\tilde{g})$ which we will denote by $\mathrm{Im}\, M_{21}^{\phi}(\tilde{g})$. Using the fact that

$$\frac{s^2\rho_1}{s^2\rho_2} \simeq \frac{m_d^2}{m_s^2} \tag{17}$$

we can safely neglect the second one such that

$$\mathrm{Im}\, M_{21}^{\varphi}(\tilde{g}) \simeq -\frac{2\alpha_s^2}{m_{\tilde{g}}^2}\,\frac{11}{54}\, R N B' c^2\theta_c\, s^2\theta_c\, s^2\rho_2\, s4\varphi\, (I_{11} - 2I_{12} + I_{22})\,. \tag{18}$$

B' represents any correction to the vacuum insertion estimate of the matrix element. The last factor appearing in the above is a manifestation of the double GIM[13] mechanism:

$$I_{11} - 2I_{12} + I_{22} \simeq \frac{1}{6\, m_{\tilde{g}}^4}\,(\tilde{M}_2^2 - \tilde{M}_1^2)^2\, \tilde{I}(\tilde{M}^2/m_{\tilde{g}}^2) \tag{19}$$

with

$$\tilde{I}(z) = \frac{1}{(1-z)^4}\left[\, 6\left(\frac{z+3}{1-z}\right)\ln z + 17 + \frac{8}{z} - \frac{1}{z^2}\right] \tag{20}$$

and $\tilde{M}$ being an average down squark mass.

Using (1), we obtain for the squark mass difference

$$\tilde{M}_2^2 - \tilde{M}_1^2 \simeq c\,(m_c^2 - m_u^2) - \rho_2^2(\mu_R^2 - \mu_L^2) \sim c\, m_c^2\,. \tag{21}$$

For simplicity we neglect $\rho_2^2\,(\mu_R^2 - \mu_L^2)$, though it could be of the order of m_c^2. With a similar omission, the appearance of the Cabibbo angle in Eq. (15) is justified [8].

From this we conclude that*

$$\mathrm{Im}\, M_{21}^{\varphi}(\tilde{g}) \simeq (1.25\times 10^{-15}\,\mathrm{GeV})\,\frac{11}{54}\, R B' c^2 \alpha_s^2\; s^2\rho_2\, s4\varphi \left(\frac{M_W}{\tilde{M}}\right)^6 x^6\, \tilde{I}(x^2) \tag{22}$$

with $x \equiv \tilde{M}/m_{\tilde{g}}$ and $m_c = 1.35$ GeV.

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* We should remark that there can be other contributions to Eq. (22) of a similar order of magnitude coming from small entries in $\Gamma_{L,R}$ which we have approximated by zero in Eq. (15), because these get multiplied by mass differences of the order $\mu_L^2 - \mu_R^2$.

To get an estimate of the phase ϕ , we now use the experimental constraint coming from the EDMN. The $\tilde{g}$ contribution to the EDM of a down quark[14] is

$$d_d(\tilde{g}) = -\frac{2e}{9\pi}\,\frac{\alpha_s}{m_{\tilde{g}}}\sum_j \mathrm{Im}\left(\Gamma_L^{jd}\Gamma_R^{jd*}\right)\frac{1}{x_j^2}\,D\left(\frac{1}{x_j^2}\right) \tag{23}$$

where D(r) is the function I_3 (r,o) defined in Ref.15). The GIM mechanism operates again and we obtain

$$\sum_j \mathrm{Im}\left(\Gamma_L^{jd}\Gamma_R^{jd*}\right)\frac{1}{x_j^2}D\left(\frac{1}{x_j^2}\right) \simeq \frac{1}{x^4}\,c\rho_1 s\rho_1 s2\varphi\left(\frac{\mu_L^2-\mu_R^2}{m_{\tilde{g}}^2}\right)\tilde{D}(x^2) \tag{24}$$

with
$$\tilde{D}(z) = \frac{z^2}{(z-1)^3}\left\{\frac{z}{2}+\frac{5}{2}-(2z+1)\frac{\ln z}{z-1}\right\}. \tag{25}$$

Using Eq (16) for $s\rho_1$ as well as the following generous upper bound

$$|d_d(\tilde{g})| \lesssim 10^{-24}\ e\cdot cm$$

we obtain
$$\alpha_s |A|\,\frac{m_{3/2}}{m_{\tilde{g}}}\left(\frac{M_W}{m_{\tilde{g}}}\right)^2\frac{1}{x^4}\,\tilde{D}(x^2)\ |s2\varphi| < 0.5\cdot 10^{-3} \tag{26}$$

which can be used as an upper bound for the phase $\phi \simeq \sin\phi$.
From Eqs.(22) and (26), we therefore conclude that

$$|\mathrm{Im}\, M_{21}^{\varphi}(\tilde{g})| \lesssim (1.25\cdot 10^{-18}\,GeV)\,\frac{11}{54}\,\frac{RB'}{|A|}\,c^2\alpha_s\,\frac{M_W}{m_{3/2}}\left(\frac{M_W}{\tilde{M}}\right)^3 s^2\rho_2\,\frac{x^6|\tilde{I}(x^2)|}{\frac{1}{x}\tilde{D}(x^2)} \tag{27}$$

where the x-dependence is rather mild in the interesting region (see Fig. 2). Taking $\alpha_S = 0.1$, $m_{3/2} = M_W$ and $s\rho_2 \sim 0.1$ as a plausible value for the mixing angle ρ_2, we conclude that we need a light down squark

$$\tilde{M} \lesssim \frac{1}{10} M_W \tag{28}$$

to get a sizeable contribution to ε_m (in order not to be in trouble with the bound on ε'/ε)[16]. Such a light $\tilde{M}$, however, is already excluded experimentallly[17].

At this point, we would like to stress that among all the gaugino contributions to M_{21} the one from the gluinos is, in fact, dominant. Indeed, the neutral $\tilde{\gamma}$ and $\tilde{Z}^0$ gaugino contributions corresponding to Figs. 1a, 1b are suppressed due to smaller gauge coupling constants. Moreover, because the functions $x^6\tilde{I}(x^2)$, $x^6\tilde{K}(x^2)$ remain finite as $x \to \infty$ (see Fig. 2), it is apparent that even a very small photino mass does not spoil this conclusion. The contribution coming from the charged gaugino $\tilde{W}$ corresponding to Fig.1c is even more suppressed[16),18].

Let us now concentrate on the contribution from the KM phase δ to Im $M_{21}(\tilde{g})$ To get an order of magnitude estimate of the LRLR term in Eq (9), we can replace the $s4\phi$ appearing in Eq. (22) by the typical KM factor $s\theta_2\, s\theta_3\, s\delta$. Even for the maximum allowed values of $s\theta_2, s\theta_3$ [19] one can easily convince oneself that this factor ($\sim 3.6\times10^{-3}$) is smaller than the maximum value of $s4\phi$ allowed by EDMN (26). Therefore its contribution is negligible. The other terms in Eq.(9) are further suppressed by a factor $s\rho_1/s\rho_2$ except for the LLLL term which, as we will show, can give an important contribution.

Now, we turn to the LLLL term. To make life easier, we ignore the LR-mixing in $M^2_{\tilde{d}}$ which is a reasonable approximation for two generations, and should provide us with the right order of magnitude estimate in the three generation case. In this situation, the coupling matrices $\Gamma_{L,R}$ are of the following simple form

$$\Gamma_L = \begin{pmatrix} K \\ 0 \end{pmatrix}, \quad \Gamma_R = \begin{pmatrix} 0 \\ 1 \end{pmatrix}. \tag{29}$$

Substituting (29) into Eq. (9) and making use of the double GIM mechanism, we obtain

$$M^{\delta}_{21}(\tilde{g}) \simeq -(1.25\cdot10^{-15}\,GeV)\; c^2 \alpha_s^2\, B\, (M_W/\tilde{M})^6 \cdot \left[1+\left(\frac{\lambda_3}{\lambda_1}\right)^2\left(\frac{m_t}{m_c}\right)^4 - 2\frac{\lambda_3}{\lambda_1}\left(\frac{m_t}{m_c}\right)^2\right]\left\{\frac{1}{27}x^6\tilde{I}(x^2)+\frac{11}{108}x^6\tilde{K}(x^2)\right\} \tag{30}$$

with

$$\lambda_i = K^*_{is} K_{id} \tag{31}$$

$$\tilde{K}(z) = \frac{2}{(1-z)^4}\left[6\left(\frac{z+1}{1-z}\right)\ln z + 10 + \frac{1}{z} + z\right] \tag{32}$$

and B being the usual correction factor to the vacuum insertion approximation in the W box diagram. We observe that Eq. (30) depends on m_t^4 [9] whereas the conventional non-SUSY box diagram has only m_t^2 dependence [11]. To estimate the imaginary part of Eq. (30), one can use the plots of $x^6\tilde{I}(x^2)$ and $x^6\tilde{K}(x^2)$ in Fig. 2 and choose values of $\tilde{M}$ greater than the experimental bound of 15 GeV [17]. It is then easy to check that for reasonable choices of λ_3 within the allowed range [19] and for even relatively light m_t ($\sim$30 GeV) one can get an ε_m of the right order.

To check that simultaneously the $\tilde{g}$-contribution to the EDMN from the KM phase δ does not exceed the experimental bound, we now replace $s2\phi$ in Eq. (24) by $s\theta_2\, s\theta_3\, s\delta$ to get a rough estimate of $d_d(\tilde{g})$ (23). As we have already noted this last factor is less than $s2\phi_{max}$. Hence the EDMN contribution from δ is safely below, but can be close to, the experimental bound, in striking contrast to the non-SUSY case[20].

Finally, we look for restrictions on $\tilde{M}$ and $m_{\tilde{g}}$ coming from the bounds on M_{21}^{δ} ($\tilde{g}$). First we consider the imaginary part. For reasonable allowed values of the top quark mass and $s\theta_2, s\theta_3, \delta$, we take [19], for example:

a) m_t = 50 GeV, $s\theta_2$ =0.05, $s\theta_3$=0.03, δ=90°

and

b) m_t = 30 GeV, $s\theta_2$ =0.07, $s\theta_3$=0.01, δ=170°

the imaginary part of the expression in the bracket in Eq. (30)

$$\mathrm{Im}\left[\left(\frac{\lambda_3}{\lambda_1}\right)^2 \left(\frac{m_t}{m_c}\right)^4 - 2\frac{\lambda_3}{\lambda_1}\left(\frac{m_t}{m_c}\right)^2 \right] \tag{33}$$

takes the values 18.2 and 0.45 respectively. Using this along with the experimental value of ε=(2.28 ± 0.05) x 10^{-3} as an upper bound for Eq. (30) and setting c=1, we obtain the plots of Fig. 3 [(curves (a) and (b)]. We also plot in Fig. 3 a more conservative bound obtained from the real part of Eq. (30) in the two generation case. [A slightly different approach has been pursued in Ref. 21 to get bounds on $\Delta\tilde{M}/\tilde{M}$ from m_K also in the four quark model]. From the curves (a) and (b), it is clear that the restrictions are rather sensitive to the top quark mass and the KM mixing angles. A definite unambiguous statement will be possible only when these are known to better accuracy. On the other hand, the restrictions from Δm_K [curve (c)], which is more reliable, turns out to be weak and does not constrain the freedom very much. In this figure, we also show the bound $\tilde{M} \gtrsim 0.9\ m_{\tilde{g}}$ coming from the absence of radiative $SU(3)_c$ and $U(1)_{EM}$ breaking through squark vacuum expectation values [22].

Conclusions:

(1) We find that supersymmetric phases alone cannot explain CP-violation in the K^0-$\bar{K}^0$ system in conjunction with the EDMN.

(2) The KM phase can give large contributions through gluino graphs implying constraints on squark and gluino masses (see Fig. 3).

(3) Moreover, gluino graphs enhance the value of the electric dipole moment of the neutron.

Finally we point out some of the limitations of our calculations: we have not considered the strong CP-violation[23)], and have also not taken into account the Lee-Weinberg mechanism[24)]. In addition, in our numerical estimations, we have always assumed the validity of the vacuum insertion method (B'=B=1) because conflicting estimates of B exist in the literature[25)].

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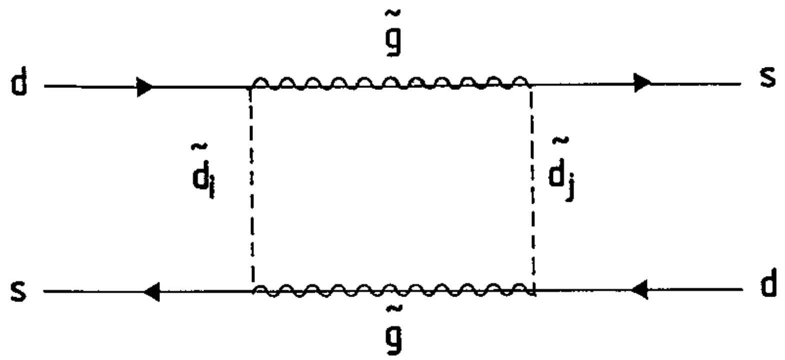
FIGURE CAPTIONS

Fig. 1 : The SUSY gluino box diagrams contributing to the $K^0\bar{K}^0$-matrix element. Fig. 1a is the usual Dirac contribution and Fig. 1b represents the additional contribution because of the Majorana nature of the gluino.

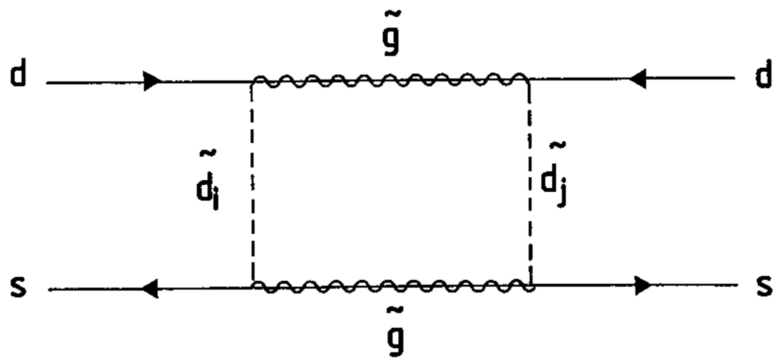
Fig. 2 : The various functions of the ratio of the squark mass to the gluino mass that enter into the CP-violation and EDMN calculations (see text).

Fig. 3 : The regions above the curves marked (a) and (b) are allowed by the CP violation parameter ε for the choices:
(a) m_t = 50 GeV, $s\theta_2$= 0.05, $s\theta_3$ = 0.03, δ =90^0
(b) m_t = 30 GeV, $s\theta_2$= 0.07, $s\theta_3$ = 0.01, δ =170^0
respectively. The broken curve (c) corresponds to the bound from $m_{K_S} - m_{K_L}$ in the two generation case. The shaded region is ruled out by experiments[17]. The region to the right of the dot-dashed line is excluded by Ref. 22).



d
s
$\tilde{g}$
$\tilde{d}_i$
$\tilde{d}_j$
s
d
$\tilde{g}$

Fig. 1a



d
d
$\tilde{g}$
$\tilde{d}_i$
$\tilde{d}_j$
s
s
$\tilde{g}$

Fig. 1b

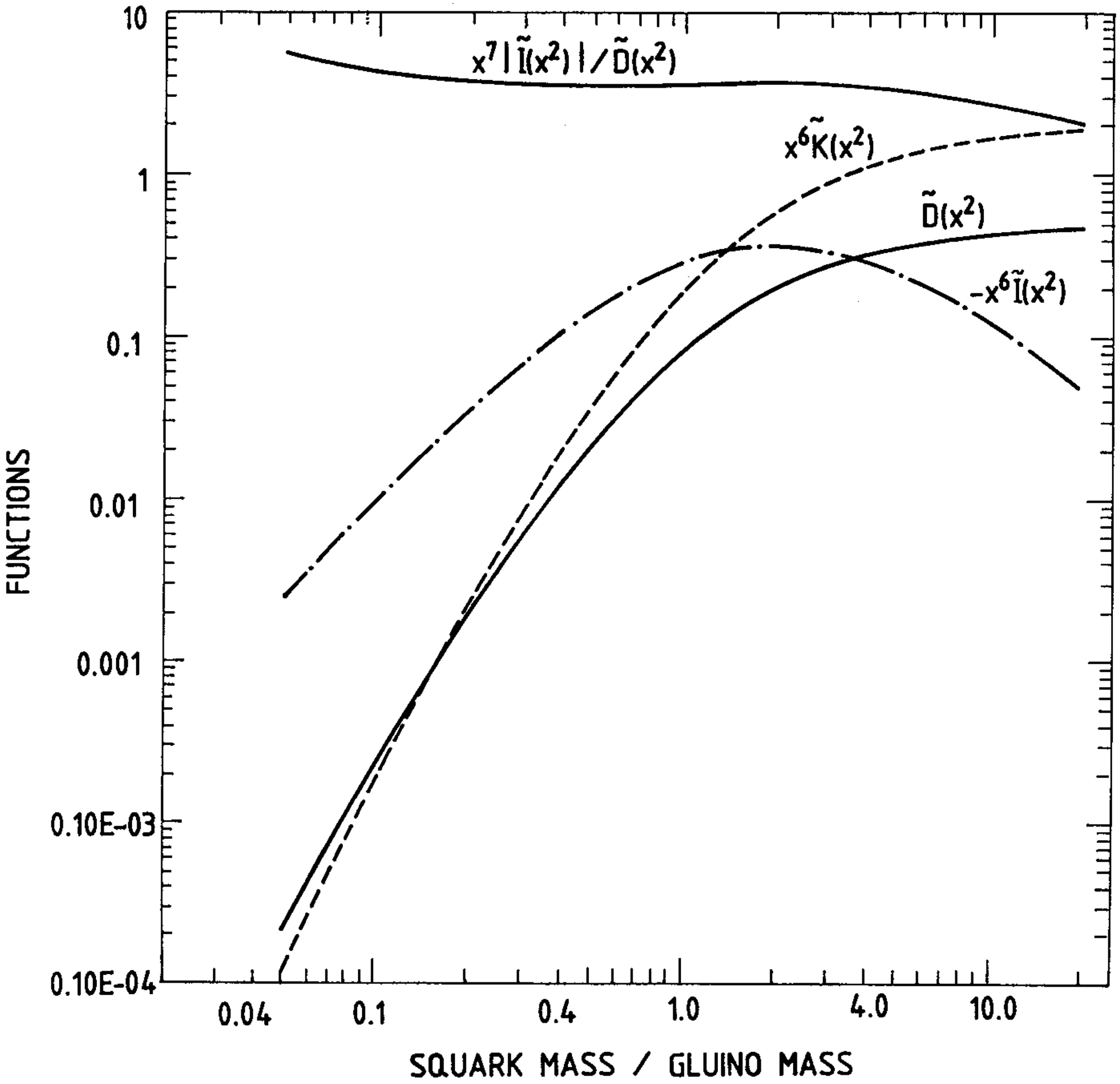

10
1
0.1
0.01
0.001
0.10E-03
0.10E-04
FUNCTIONS
0.04
0.1
0.4
1.0
4.0
10.0
SQUARK MASS / GLUINO MASS
x⁷|Ĩ(x²)|/D̃(x²)
x⁶K̃(x²)
D̃(x²)
-x⁶Ĩ(x²)

Fig. 2

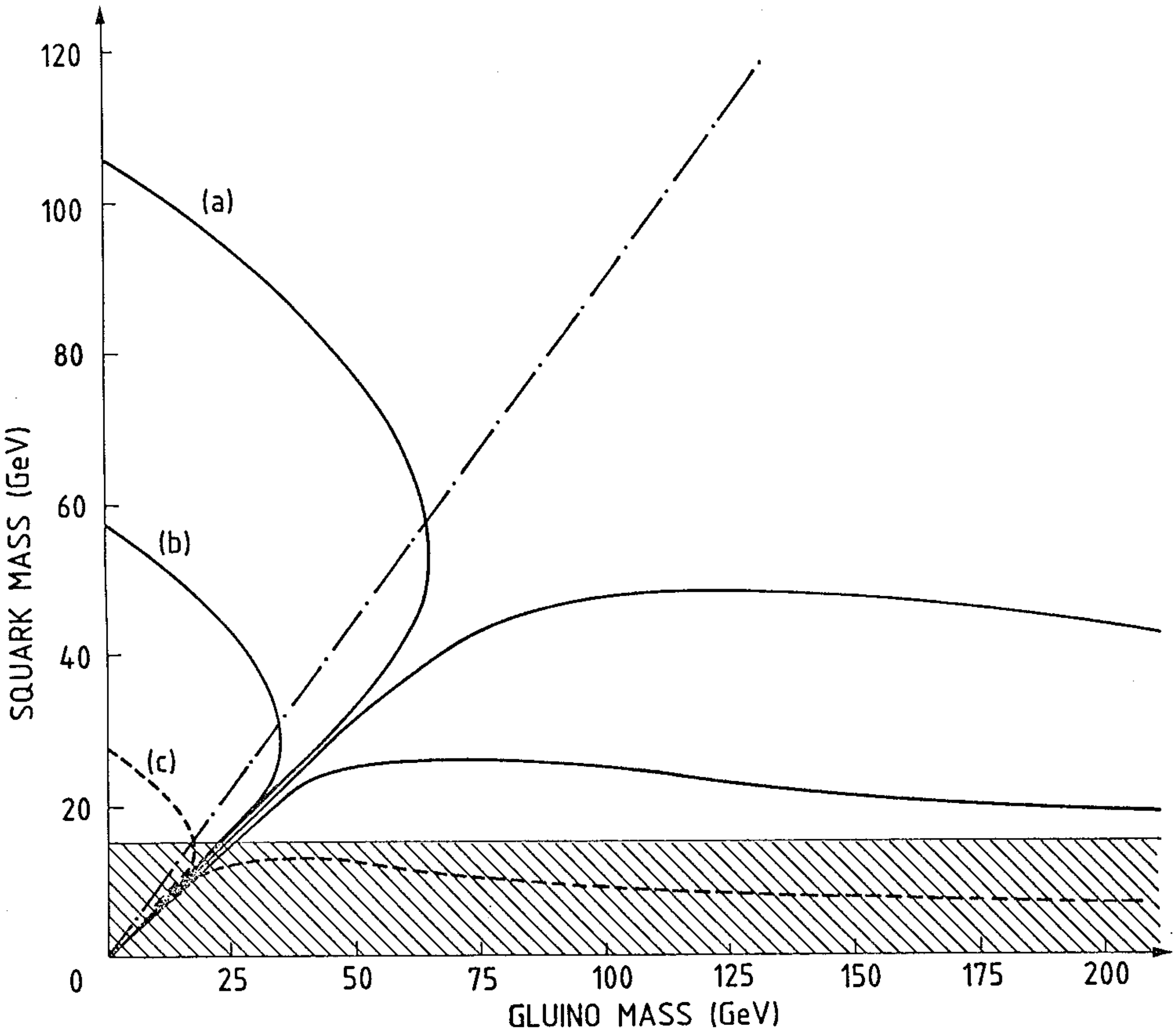

SQUARK MASS (GeV)
120
100
80
60
40
20
0
(a)
(b)
(c)
25
50
75
100
125
150
175
200
GLUINO MASS (GeV)

Fig. 3