Ref.TH.3927-CERN

GLUINO PENGUINS AND ε'/ε

J.-M. Gérard*, W. Grimus+ and Amitava Raychaudhuri†

CERN - Geneva

ABSTRACT

We calculate the penguin diagrams involving gluinos and discuss their implications on ε'/ε. The dominant contribution originates in the Kobayashi-Maskawa phase. Fixing this phase from the ε parameter, we find a smaller value of ε'/ε than in the non-supersymmetric case alone.

---

*) IISN fellow, on leave of absence from Institut de Physique Théorique, Université de Louvain, Louvain-La-Neuve (Belgium); Address after September 1984, Max-Planck-Institut für Physik and Astrophysik, Munich (Germany).

+) Address from July 1, Institut für theoretische Physik, Boltzmanngasse 5, Vienna (Austria).

†) On leave (1983-84) from the Department of Pure Physics, Calcutta University, Calcutta (India).

Ref.TH.3927-CERN
June 1984

Ever since its experimental discovery in the Kaon system, CP-violation has been a fascinating field of theoretical research. Unfortunately, up until now no CP-violating effect has been found elsewhere, so the main focus is still on K^0-$\bar{K}^0$. In this paper we will concentrate on the supersymmetry (SUSY) contributions to the CP-violation parameter ε'/ε which enters in the measured amplitude ratios [1] η_{+-}, η_{00}:

$$\eta_{+-} = \frac{A(K_L \to \pi^+\pi^-)}{A(K_S \to \pi^+\pi^-)} = \varepsilon + \frac{\varepsilon'}{1+\omega/\sqrt{2}}$$

$$\eta_{00} = \frac{A(K_L \to \pi^0\pi^0)}{A(K_S \to \pi^0\pi^0)} = \varepsilon - \frac{2\varepsilon'}{1-\omega\sqrt{2}}$$

$$\omega = e^{i(\delta_2-\delta_0)} \mathrm{Re}\, A_2 / \mathrm{Re}\, A_0 \qquad (|\omega| \approx 1/20) \tag{1}$$

where A_0, A_2 are the I=0,2 amplitudes for $K^0 \to \pi\pi$ and δ_0, δ_2 are the strong interaction phase shifts. These equations are first order relations in the CP-violating quantities $\varepsilon, \varepsilon'$. So far, the question if $\eta_{+-} = \eta_{00}$ i.e., $\varepsilon'=0$ has not been settled; there exists only a bound [2]:

$$\frac{\varepsilon'}{\varepsilon} \doteq -0.0046 \pm 0.0053 \pm 0.0024 \tag{2}$$

but promising experiments are being undertaken or planned. This question is important because the very successful standard model which also offers the nice possibility of explaining CP-violation through the Kobayashi-Maskawa (KM) mechanism also predicts a non-vanishing ε'. Recently however it has been claimed that the standard model might not be able to account for ε if the top quark turned out to be light [3]. The reason is that the KM angles θ_2, θ_3 are constrained to be rather small [4] because of the relatively long b-lifetime [5] and the small ratio $\overline{R} = \Gamma(b\to u e\nu)/\Gamma(b\to c e\nu) \lesssim 0.04$ [6].

Writing ε and ε' in more detail one obtains:

$$\varepsilon = \frac{e^{i\frac{\pi}{4}}}{2\sqrt{2}}\left(\varepsilon_m + 2\xi\right)$$

$$\varepsilon' = \frac{i\omega}{\sqrt{2}}\left(\frac{\mathrm{Im}\,A_2}{\mathrm{Re}\,A_2} - \xi\right) \simeq -i\omega\,\xi/\sqrt{2}$$

$$\varepsilon_m = \frac{\mathrm{Im}\,M_{12}}{\mathrm{Re}\,M_{12}}\,,\quad \xi = \frac{\mathrm{Im}\,A_0}{\mathrm{Re}\,A_0} \tag{3}$$

where ξ is given by the so-called penguin diagrams [7,8]. In the standard model Im A_2/Re A_2 is very small and we will neglect it. There are two major theoretical uncertainties in the calculation of the K^0-$\overline{K}^0$ transition matrix element M_{12}. Firstly, long distance contributions [9] to M_{12} cannot be calculated reliably and secondly, in the box diagram calculation [10] there appears an inaccurately known parameter B which is the correction factor to the matrix element calculated by vacuum insertion. In general long distance contributions can modify the CP predictions of the box diagram considerably [11]; however, arguments have been advanced that there may be important cancellations in the imaginary part of M_{12} [12]. Both large B and long distance effects in M_{12} could conspire in such a way as to give enough CP-violation in the standard model. On the other hand, it has been shown that even in the worst case for the standard model, small B [13]

(B=0.33), low top quark mass, m_t, and no long distance contributions, CP-violation can be satisfactorily explained [14,15] by going to the supersymmetric extension and including soft breaking terms [16]. In particular, the gluino ($\tilde{g}$) - squark ($\tilde{q}$) box diagram can account for the missing CP-violation in M_{12} where the relevant source of CP-violation in the $\tilde{g}$-$\tilde{d}$-d interaction is again the KM phase δ. A new phase ϕ, the "SUSY phase", can only enter in contributions to M_{12} through $\tilde{d}_L$ - $\tilde{d}_R$ mixing but its effect on ε is very small. This phase is a combination of possible phases in the gluino mass ($\tilde{m}$) and in the parameter A which appears in the soft SUSY breaking trilinear scalar couplings [16]. In view of this interesting result, it is important to also examine the SUSY contribution to ε', i.e., the superpenguins [17]. In this paper we make a detailed study of the penguins involving gluinos, the 'pengluinos' (Fig. 1), which constitute the most important SUSY contribution to ε' originating in the KM phase δ. Using restrictions for the SUSY phase ϕ from the Electric Dipole Moment of the Neutron (EDMN) [18,14], we also look for additional effects of ϕ in ε'.

We begin by writing down the $\tilde{g}$-$\tilde{d}$-d interaction:

$$\mathcal{L}_{\tilde{g}\tilde{d}d} = i\sqrt{2}\, g_s\, \tilde{d}_i^{\dagger}\, \bar{\tilde{g}}_a \frac{\lambda^a}{2} \left(\Gamma_L P_L + \Gamma_R P_R \right)^{ij} d_j + h.c.$$

$$P_L = \frac{1-\gamma_5}{2}, \quad P_R = \frac{1+\gamma_5}{2} \tag{4}$$

where g_s denotes the coupling constant of the strong interactions and the λ^a's are the Gell-Mann matrices of $SU(3)_c$. In the basis where the quark mass matrices are diagonal ($\hat{M}_{u,d}$), the down squark mass matrix is

$$M^2_{\tilde{d}} = \begin{pmatrix} \mu_L^2 1 + \hat{M}_d^2 + c K^\dagger \hat{M}_u^2 K & |A| \mu \hat{M}_d \\ |A| \mu \hat{M}_d & \mu_R^2 1 + \hat{M}_d^2 \end{pmatrix} \tag{5}$$

K is the KM matrix, $\mu_{L,R}$ and μ come from soft SUSY breaking terms. The parameter 'c' characterises radiative corrections to the tree level squark mass matrix which introduce flavour violation [19]. It is in general between 0.1 and 1 and negative. In most of our calculations we will take $c = -1$, which favours SUSY contribution to ε. The 6x3 matrices Γ_L, Γ_R appearing in Eq. (4) are given by [14]

$$(\Gamma_L, \Gamma_R) = U^\dagger \begin{pmatrix} e^{i\phi} 1 & 0 \\ 0 & e^{-i\phi} 1 \end{pmatrix} \tag{6}$$

where U diagonalises the down squark mass matrix in (5). The mechanism of flavour violation through Eq. (4) explained above also leads to interesting results in the B^0-$\bar{B}^0$ system, i.e. large mixing and enhanced directly measurable CP-violating quantities [20].

To establish our conventions we exhibit in Fig. 2 the necessary Feynman rules for the calculation of the pengluinos. After having done the momentum integration, we obtain for the graphs of Fig. 1

$$(D_a)^\mu_\alpha = -\frac{i}{16\pi^2} g_s^3 C_2(G)\frac{\lambda^a}{2}\bar{O}_j \int_0^1 dx \int_0^{1-x} dy \left\{ \left(\frac{1}{2-\frac{n}{2}} - \gamma - 1 - \ln\frac{\bar{D}_j}{4\pi\mu^2}\right)\gamma^\mu + \right.$$

$$\left. + (x\not{p} + (1-y)\not{k} + \tilde{m})\gamma^\mu (x\not{p} - y\not{k} + \tilde{m}) / \bar{D}_j \right\} O_j$$

$$\bar{D}_j = M_j^2 x + \tilde{m}^2(1-x) - x(1-x)p^2 - y(1-y)k^2 - 2k\cdot p\, xy$$

$$(D_b)^\mu_\alpha = -\frac{i}{16\pi^2} g_s^3 (2C_2(R) - C_2(G))\frac{\lambda^a}{2}\bar{O}_j \int_0^1 dx \int_0^{1-x} dy \left\{ \left(\frac{1}{2-\frac{n}{2}} - \gamma - \right.\right.$$

$$\left. - \ln\frac{D_j}{4\pi\mu^2}\right)\gamma^\mu + \left((1-x)\not{p} + y\not{k} + \tilde{m}\right)(2xp + (1-2y)k)^\mu / D_j \Big\} O_j$$

$$D_j = \tilde{m}^2 x + M_j^2(1-x) - x(1-x)p^2 - y(1-y)k^2 - 2k\cdot p\, xy$$

$$(D_c + D_d)^\mu_\alpha = -ig_s\frac{\lambda^a}{2}\left(\gamma^\mu \frac{1}{\not{p} - m_d}\Sigma(p) + \Sigma(p')\frac{1}{\not{p}' - m_s}\gamma^\mu\right)$$

$$-i\Sigma(p) = \frac{i}{16\pi^2}\cdot 2g_s^2 C_2(R)\bar{O}_j \int_0^1 dx \left((1-x)\not{p} + \tilde{m}\right)\left(\frac{1}{2-\frac{n}{2}} - \gamma - \ln\frac{E_j}{4\pi\mu^2}\right) O_j$$

$$E_j = (1-x)M_j^2 + x\tilde{m}^2 - x(1-x)p^2, \quad k = p' - p \; . \tag{7}$$

M_j denotes the down squark masses. We have used dimensional regularisation (n = # of dimensions, γ = Euler constant and μ is a scale parameter with the dimension of mass). O_j, $\bar{O}_j$ are given by

$$O_j = \Gamma_L^{js} P_L + \Gamma_R^{js} P_R \,, \quad \bar{O}_j = \Gamma_L^{jd*} P_R + \Gamma_R^{jd*} P_L \,. \tag{8}$$

In order to keep track of the origin of the different terms in Eq. (7), we have retained the quadratic Casimir operators $C_2(R)$, $C_2(G)$ of the fundamental and the adjoint representations respectively. For SU(N) their eigenvalues are:

$$C_2(R) = \frac{N^2-1}{2N} \,, \quad C_2(G) = N \tag{9}$$

In Eq. (7), the infinities proportional to $C_2(G)$ cancel each other as they must. The reason is the conservation of the gluonic current ($SU(3)_c$ is unbroken) and the Ward - Takahashi identity

$$\begin{aligned} k^\mu \Gamma_\mu^a &= g_s \frac{\lambda^a}{2} \left(\Sigma(p) - \Sigma(p') \right) \\ \Gamma_\mu^a &= i \left(D_a + D_b \right)_\mu^a \end{aligned} \tag{10}$$

which is preserved by the renormalised quantities. Because the self-energy Σ contains only $C_2(R)$ the infinities proportional to $C_2(G)$ must cancel each other in $\Gamma_\mu{}^a$. But, of course, as the alert reader has already noticed, in this case all infinities are cancelled anyway by the GIM mechanism:

$$\Gamma_A^{jd*} \Gamma_B^{js} = 0 \,; \quad A, B = L, R \tag{11}$$

However, the previous observations are independent of (11) and serve as a check. Finally, we want to remark that on mass shell, using the renormalised vertex $\Gamma_{R\mu}{}^a$ which is obtained by the renormalisation conditions (10) and [21]

$$\Sigma_R(p)\,u_s(p)=0\,,\quad \bar u_d(p')\,\Sigma_R(p')=0 \tag{12}$$

(Σ, Σ_R denote the transition self-energies s→d) is equivalent to summing over all (unrenormalised) diagrams a,b,c,d. Taking the latter option, after having put the wings of the pengluino on shell, one obtains expressions corresponding to Eq. (7)

$$\bar u_d(p')\,i\,(D_a)^a_\mu\,u_s(p)=\frac{g_s^3}{16\pi^2}C_2(G)\,\bar u_d(p')\frac{\lambda^a}{2}\int_0^1 dx\int_0^{1-x}dy\Big\{\Big[\gamma_\mu\Big(-\ln\frac{\bar D_j}{\tilde m^2}+$$

$$+\frac{1}{\bar D_j}\big((y-y^2-xy)k^2+\tilde m^2\big)\Big)+\frac{1}{\bar D_j}xy\,m_d\,i\sigma_{\mu\nu}k^\nu+\frac{1}{\bar D_j}k_\mu m_d(xy-2y+2y^2)\Big]\cdot$$

$$\cdot\left(\Gamma_L^{jd*}\Gamma_L^{js}P_L+\Gamma_R^{jd*}\Gamma_R^{js}P_R\right)+\frac{1}{\bar D_j}\Big[\gamma_\mu m_d m_s x^2+m_s(x-xy-x^2)\,i\sigma_{\mu\nu}k^\nu+$$

$$+k_\mu m_s\,(2y-2y^2-3xy+x-x^2)\Big]\left(\Gamma_L^{jd*}\Gamma_L^{js}P_R+\Gamma_R^{jd*}\Gamma_R^{js}P_L\right)+$$

$$+\frac{1}{\bar D_j}\Big[\gamma_\mu\tilde m\,m_d\,x+\tilde m(1-x)\,i\sigma_{\mu\nu}k^\nu+k_\mu\tilde m\,(1-x-2y)\Big]\Big(\Gamma_L^{jd*}\Gamma_R^{js}P_R+$$

$$+\Gamma_R^{jd*}\Gamma_L^{js}P_L\Big)+\frac{1}{\bar D_j}\gamma_\mu\,\tilde m\,m_s\,x\left(\Gamma_L^{jd*}\Gamma_R^{js}P_L+\Gamma_R^{jd*}\Gamma_L^{js}P_R\right)\Big\}\,u_s(p)$$

$$\bar u_d(p')\,i\,(D_b)^a_\mu\,u_s(p)=\frac{g_s^3}{16\pi^2}\big(2C_2(R)-C_2(G)\big)\bar u_d(p')\frac{\lambda^a}{2}\int_0^1 dx\int_0^{1-x}dy\;\cdot$$

$$\cdot\Big\{\Big[\gamma_\mu\Big(-\ln\frac{D_j}{\tilde m^2}+\frac{1}{D_j}\big(xy\,m_d^2+x(1-x-y)\,m_s^2\big)\Big)-\frac{1}{D_j}xy\,m_d\,i\sigma_{\mu\nu}k^\nu+$$

$$+ \frac{1}{D_j} m_d k_\mu (1-2y-x) y \Big] \left(\Gamma_L^{jd*} \Gamma_L^{js} P_L + \Gamma_R^{jd*} \Gamma_R^{js} P_R \right) +$$

$$+ \frac{1}{D_j} \Big[\gamma_\mu (1-x) x\, m_d m_s - m_s x (1-x-y) i \sigma_{\mu\nu} k^\nu + m_s k_\mu (1-x-y)(1-x-2y) \Big] \cdot$$

$$\cdot \left(\Gamma_L^{jd*} \Gamma_L^{js} P_R + \Gamma_R^{jd*} \Gamma_R^{js} P_L \right) + \frac{1}{D_j} \Big[\gamma_\mu \tilde{m} m_d x - \tilde{m} x i \sigma_{\mu\nu} k^\nu +$$

$$+ \tilde{m} k_\mu (1-2y-x) \Big] \left(\Gamma_L^{jd*} \Gamma_R^{js} P_R + \Gamma_R^{jd*} \Gamma_L^{js} P_L \right) +$$

$$+ \frac{1}{D_j} \gamma_\mu \tilde{m} m_s x \left(\Gamma_L^{jd*} \Gamma_R^{js} P_L + \Gamma_R^{jd*} \Gamma_L^{js} P_R \right) \Big\} u_s(p)$$

$$\bar{u}_d(p') \, i (D_c + D_d)^a_\mu u_s(p) = \frac{g_s^3}{16\pi^2} \cdot 2 C_2(R) \bar{u}_d(p') \frac{\lambda^a}{2} \gamma_\mu \cdot$$

$$\cdot \Big\{ \frac{f_{1j}(m_s^2) m_s^2 - f_{1j}(m_d^2) m_d^2}{m_s^2 - m_d^2} \left(\Gamma_L^{jd*} \Gamma_L^{js} P_L + \Gamma_R^{jd*} \Gamma_R^{js} P_R \right) +$$

$$+ m_s m_d \frac{f_{1j}(m_s^2) - f_{1j}(m_d^2)}{m_s^2 - m_d^2} \left(\Gamma_L^{jd*} \Gamma_L^{js} P_R + \Gamma_R^{jd*} \Gamma_R^{js} P_L \right) +$$

$$+ \frac{f_{2j}(m_s^2) - f_{2j}(m_d^2)}{m_s^2 - m_d^2} \Big[\tilde{m} m_d \left(\Gamma_L^{jd*} \Gamma_R^{js} P_R + \Gamma_R^{jd*} \Gamma_L^{js} P_L \right) +$$

$$+ \tilde{m} m_s \left(\Gamma_L^{jd*} \Gamma_R^{js} P_L + \Gamma_R^{jd*} \Gamma_L^{js} P_R \right) \Big] \Big\} u_s(p)$$

$$f_{1j}(p^2) = \int_0^1 dx\, (1-x) \ln \frac{E_j}{\tilde{m}^2}, \quad f_{2j}(p^2) = \int_0^1 dx \ln \frac{E_j}{\tilde{m}^2} \qquad (13)$$

Here m_s, m_d should be considered as constituent quark masses. Since squarks and gluinos are expected to be much heavier than the ordinary hadronic scale, we have

$$\tilde{m}^2 x + M_j^2 (1-x) \gg m_s^2, m_d^2, k^2 \tag{14}$$

for all $x \in [0,1]$ and we can safely make an expansion in the latter quantities which are all of the same order and neglect the terms of higher order. We thus obtain our final result for the pengluinos:

$$\bar{u}_d(p')\, i\,(D_a+D_b+D_c+D_d)^a_\mu\, u_s(p) = \frac{g_s^3}{16\pi^2\tilde{m}^2}\, \bar{u}_d(p')\frac{\lambda^a}{2}\cdot$$

$$\cdot\Big\{\tfrac{1}{3}(k^2\gamma_\mu - k_\mu \not{k})(C_2(G)A(z_j)+C_2(R)B(z_j))(\Gamma_L^{jd*}\Gamma_L^{js}P_L+\Gamma_R^{jd*}\Gamma_R^{js}P_R)+$$

$$+\, i\sigma_{\mu\nu}k^\nu\Big[\big(C_2(G)C(z_j)-C_2(R)D(z_j)\big)\Big(m_d\,(\Gamma_L^{jd*}\Gamma_L^{js}P_L +$$

$$+\,\Gamma_R^{jd*}\Gamma_R^{js}P_R) + m_s\,(\Gamma_L^{jd*}\Gamma_L^{js}P_R+\Gamma_R^{jd*}\Gamma_R^{js}P_L)\Big) +$$

$$+\,\big(C_2(G)E(z_j)-4C_2(R)C(z_j)\big)\,\tilde{m}\,(\Gamma_L^{jd*}\Gamma_R^{js}P_R+\Gamma_R^{jd*}\Gamma_L^{js}P_L)\Big]\Big\}u_s(p)$$

$$z_j = \frac{M_j^2}{\tilde{m}^2} \tag{15}$$

The functions A,B,C,D,E are given by

$$A(z) = \frac{1}{2(1-z)^2}\left\{3-3z+(2z-1)\ln z\right\}$$

$$B(z) = \frac{1}{(1-z)^4}\left\{-\frac{11}{6}+3z-\frac{3}{2}z^2+\frac{1}{3}z^3-\ln z\right\}$$

$$C(z) = \frac{1}{4(1-z)^3}\left\{1-z^2+2z\ln z\right\}$$

$$D(z) = \frac{1}{(1-z)^4}\left\{\frac{1}{3}+\frac{1}{2}z-z^2+\frac{1}{6}z^3+z\ln z\right\}$$

$$E(z) = \frac{1}{(1-z)^2}\left\{1-z+z\ln z\right\} \qquad (16)$$

In the ordinary penguin there is no $\sigma_{\mu\nu}k^\nu$-term which mixes left- and right-handed quarks. Moreover, the $m_{s,d}\sigma_{\mu\nu}k^\nu$-terms which can in principle be of the same order as the $\gamma_\mu k^2$-term are suppressed even if m_t is large [21,22]. In the SUSY case we have no reason to expect a suppression of the $\sigma_{\mu\nu}k^\nu$-terms and this is borne out by our numerical results.

To calculate the phase of the amplitude $A_0(K^0\to\pi\pi)$ we have to estimate the matrix elements of the relevant operators. One of them is the well-known operator [23]

$$O_5 = \bar{d}_L\gamma_\mu\lambda^a s_L\left(\bar{u}_R\gamma^\mu\lambda^a u_R+\bar{d}_R\gamma^\mu\lambda^a d_R+\bar{s}_R\gamma^\mu\lambda^a s_R\right) \qquad (17)$$

where the gluon propagator is cancelled by the k^2 from the vertex whereas the other one containing $\sigma_{\mu\nu}k^\nu$ is non-local. To evaluate the matrix elements we use the vacuum insertion method [23] with the modification proposed in Ref. [24] to also include momentum dependence in the pseudoscalar amplitudes, e.g.

$$\langle 0|\bar{d}\gamma_5 s|\bar{K}^0(k)\rangle = -i\sqrt{2}f_K \frac{m_K^2}{m_d+m_s}\frac{1}{1-k^2/m_\sigma^2} \tag{18}$$

in order to be consistent with PCAC ($f_K \approx f_\pi = 93$ MeV, m_σ is the mass of the 0^+ scalar meson and here the m_d, m_s should be considered as current quark masses). To evaluate the matrix element of the second operator [25] we replace the gluon propagator by M_H^{-2}, where M_H is a typical hadronic mass scale, a procedure which should at least give us the right order of magnitude. In the numerical analysis we will choose M_H to be the Kaon mass. Now we can write down our results for

$$\xi = -f\, \mathrm{Im}\, P/\mathrm{Re}\, P \tag{19}$$

where

$$P = \frac{G_F}{\pi\sqrt{2}}\, G_A - \frac{\alpha_s}{3\tilde{m}^2}\left(C_2(G)A(z_j)+C_2(R)B(z_j)\right)\left(\Gamma_L^{jd*}\Gamma_L^{js}+\Gamma_R^{jd*}\Gamma_R^{js}\right)-$$

$$-\frac{\alpha_s}{\tilde{m}^2}\Big\{\left(C_2(G)C(z_j)-C_2(R)D(z_j)\right)\left(\Gamma_L^{jd*}\Gamma_L^{js}-\Gamma_R^{jd*}\Gamma_R^{js}\right)(m_s-m_Q)+$$

$$+\left(C_2(G)E(z_j)-4C_2(R)C(z_j)\right)\left(\Gamma_L^{jd*}\Gamma_R^{js}-\Gamma_R^{jd*}\Gamma_L^{js}\right)\tilde{m}\Big\}.$$

$$\cdot \frac{3m_Q}{M_H^2} \frac{f_+ + \frac{f_K}{f_\pi}\left(1+2\frac{m_K^2}{m_\sigma^2}\right)}{-f_+ + \frac{f_K}{f_\pi}\left(1+2\frac{m_K^2}{m_\sigma^2}\right)} \tag{20}$$

with

$$G_A = \frac{2}{3}\left(\lambda_c \ln\frac{m_c^2}{\mu^2} + \lambda_t \ln\frac{m_t^2}{\mu^2}\right), \quad \lambda_j = K_{jd}^* K_{js} \,. \tag{21}$$

G_A comes from the contribution of the ordinary penguin and is valid in the limit $m_t \ll M_W$ (in our calculation we use the exact expression of Ref. [21]). In this equation μ is again a typical hadronic mass scale related to the momentum transfer involved. As before, we will take μ to be the Kaon mass for our numerics. f is defined as the fraction of the total amplitude A_0 which can be attributed to penguins [8]. α_s is the strong coupling constant which appears in the loop and will be fixed at 0.1. m_Q denotes the u,d constituent quark mass; we will choose m_Q = 300 MeV and m_s = 500 Mev. f_+ is defined by

$$\langle \pi^+(q) | \bar{u}\gamma_\mu s | \bar{K}^0(k) \rangle = f_+(k+q)_\mu + f_-(k-q)_\mu \tag{22}$$

and its numerical value is $f_+ \approx 1$. Since in evaluating the necessary matrix elements we neglect $m_\pi{}^2/m_K{}^2$, f_- does not appear in (20). Finally, for m_σ we will take 700 Mev.

As far as the ΔI=1/2 rule is concerned [23], penguin contributions are negligible for reasonable values of the squark and gluino masses [26] (in the numerical examples below, the real part of the SUSY penguins does not exceed 5% of that of the conventional penguin). It is therefore legitimate to take the value of the parameter f from non-SUSY considerations. We choose f=1/6 to be in agreement with Ref. [27].

Before turning to the discussion of the numerical results, we briefly describe our procedure [20,3]. Given the bottom lifetime τ_B and the ratio $\bar{R}$, the KM angle θ_3 is fixed and θ_2 is obtained as a function of the KM phase δ. Now, for a fixed m_t, the experimental value of ε is used to fix δ, where the theoretical expression for ε is given by the sum of the ordinary W-box and the gluino box diagrams. The contribution of ξ is neglected because of the smallness of ε'/ε (see numerical results below). As in the standard case [27], for every τ_B there exist two solutions for δ for both of which $\sin\delta > 0$. We will vary τ_B but will always keep $\bar{R}=0.03$.

Let us now mention the SUSY input values and the range of the top quark mass. To get a significant SUSY contribution to ε we choose a large c (c=-1) and rather light SUSY mass parameters, namely $\tilde{m}$=40 GeV, μ_L = 50 GeV, $\mu_L^2 - \mu_R^2$ = 100 GeV² and μ = 40 GeV. As it can be easily seen from $M^2_{\tilde{d}}$, one obtains an upper bound on m_t

$$m_t \lesssim \mu_L / \sqrt{-c} \tag{23}$$

in order to keep $SU(3)_c \times U(1)_{em}$ unbroken (i.e. all $M_j^2 > 0$). This is a consequence of $c < 0$. Therefore, in this case, $m_t \lesssim 45$ GeV to have a reasonably heavy $\tilde{b}_L$, which is the lightest down squark in this scenario. Finally, we should stress again that we always take B=0.33.

As far as our numerical results are concerned, one should keep in mind that there is some uncertainty involved in the ratio of the matrix elements in Eq. (20) which might affect the relative size of the γ_μ and $\sigma_{\mu\nu}k^\nu$-contributions; nonetheless the general features should remain unchanged.

It turns out that ε'/ε is somewhat diminished by the SUSY contribution. This point deserves some elaboration. The SUSY contribution to ε'/ε consists of three terms which correspond to γ_μ, $m_{s,d}\sigma_{\mu\nu}k^\nu$ and $\tilde{m}\sigma_{\mu\nu}k^\nu$ in Eq. (15) and are of the same order as the ordinary contribution. It can be checked that for negative c, the last one, which is proportional to $\tilde{d}_L$-$\tilde{d}_R$ mixing, has the same sign as the ordinary contribution, while the other two are of opposite sign. Therefore for A=0, the SUSY contribution to ε'/ε is negative, whereas for $|A|=3$, it turns out that the last term dominates and the net SUSY contribution is positive. However, the total (ordinary + SUSY) ε'/ε is always less than the value obtained from the non-SUSY theory alone. This is due to the fact that when SUSY contributions are included, the ε parameter can be fitted with a somewhat smaller KM phase δ and even the ordinary contribution to ε'/ε is thereby reduced.

The SUSY phase ϕ can only have an effect through the $\tilde{m}\sigma_{\mu\nu}k^\nu$-term and for $A\neq 0$. But the numerical calculations show that with the bound on ϕ from the EDMN ($|\phi| \lesssim 10^{-2}$), it has at most an effect of 10% on ε'/ε. Since this would shift the curves only marginally we do not show it in the plots.

In Fig. 3 we exhibit the behaviour of ε'/ε as a function of m_t. We have taken $|A|=3$ and have plotted curves for several values of the bottom lifetime τ_B. Since ε'/ε is a double- valued function of τ_B, it must also be so as a function of m_t. This is indeed so in Fig. 3 where for comparison we have also presented the ordinary curve for $\tau_B=0.5\times10^{-12}$sec; for the range of m_t in this figure, the other two values of τ_B are excluded in the ordinary case by the ε parameter. In Fig. 4 the top quark mass is fixed at 40 GeV and the behaviour as a function of τ_B for different choices of $|A|$ and c is shown. For A=0 there is a remarkable, although accidental, cancellation between the SUSY and the ordinary contributions and ε'/ε is even slightly negative for the input values we have taken. In some situations therefore Super-Kobayashi-Maskawa CP-violation mimics the Superweak model predictions. Finally, we should point out that we have used $\bar{R} = 0.03$. For smaller $\bar{R}$, the curves lie within the corresponding ones that we have presented.

In conclusion, we can say that in addition to the nice features of CP-violation induced by gluinos in K^0-$\bar{K}^0$ mixing, this mechanism is not endangered by EDMN [20] and ε'/ε. The prediction for ε'/ε is smaller than that in the standard case for negative values of the flavour violation parameter c, which is preferred by radiative corrections. Furthermore, our calculations indicate that $\tilde{d}_L$-$\tilde{d}_R$ mixing is of significance as far as ε'/ε is concerned.

After we had finished this work, we received a preprint by P. Langacker and B. Sathiapalan (Univ. of Pennsylvania preprint UPR-0256T) which also discusses the gluino contribution to ε'/ε. Our study is however more complete because we include $\tilde{d}_L$-$\tilde{d}_R$ mixing and the $m_{s,d}\sigma_{\mu\nu}k^\nu$ term.

ACKNOWLEDGEMENTS

It is a great pleasure to thank A. Masiero and D.V. Nanopoulos for their interest in this work. W.G. enjoyed many informative discussions with G. Ecker on penguins and CP-violation.

References

[1] T.D. Lee and C.S. Wu, Ann. Rev. Nucl. Science 16 (1966) 511;
L. Wolfenstein, in Theory and Phenomenology in Particle Physics, Ed. A. Zichichi, Academic Press, New York (1969).

[2] Univ. of Chicago - Saclay Collaboration: results presented by B. Winstein at C.E.R.N. (June 1984).

[3] P.H. Ginsparg, S.L. Glashow and M.B. Wise, Phys. Rev. Lett. 50 (1983) 1415;
A.J. Buras, W. Slominski and H. Steger, Nucl. Phys. B238 (1984) 529.

[4] K. Kleinknecht and B. Renk, Phys. Lett. 130B (1983) 459.

[5] E. Fernandez et. al. (MAC Collaboration), Phys. Rev. Lett. 51 (1983) 1022;
N.S. Lockyer et. al. (Mark II Collaboration), Phys. Rev. Lett. 51 (1983) 1316.

[6] A. Chen et. al. (CLEO Collaboration), Phys. Rev. Lett. 52 (1984) 1084.

[7] J. Ellis, M.K. Gaillard and D.V. Nanopoulos, Nucl. Phys. B109 (1976) 213.

[8] F.J. Gilman and M.B. Wise, Phys. Lett. 83B (1979) 83.

[9] L. Wolfenstein, Nucl. Phys. B160 (1979) 501.

[10] M.K. Gaillard and B.W. Lee, Phys. Rev. D10 (1974) 897.

[11] J.F. Donoghue and B.R. Holstein, Phys. Rev. D29 (1984) 2088;
G. Ecker, Univ. of Vienna preprint UWThPh-1984-21(1984).

[12] C.T. Hill, Phys. Lett. 97B (1980) 275.

[13] J. Donoghue, E. Golowich and B. Holstein, Phys. Lett. 119B (1982) 412.

[14] J.-M. Gérard, W. Grimus, Amitava Raychaudhuri and G. Zoupanos, Phys. Lett. 140B (1984) 349.

[15] J.-M. Gérard, W. Grimus, A. Masiero, D.V. Nanopoulos and Amitava Raychaudhuri, CERN preprint TH.3837 (to be published in Phys. Lett. B).

[16] A.H. Chamseddine, R. Arnowitt and P. Nath, Phys. Rev. Lett. 49 (1982) 970;
R. Barbieri, S. Ferrara and C.A. Savoy, Phys. Lett. 119B (1982) 343;
H.P. Nilles, M. Srednicki and D. Wyler, Phys. Lett. 120B (1982) 346;
L. Hall, J. Lykken and S. Weinberg, Phys. Rev. D27 (1983) 2359;
S.K. Soni and H.A. Weldon, Phys. Lett. 126B (1983) 215.

[17] S.-P. Chia and S. Nandi, Phys. Rev. D27 (1983) 1654;
E. Franco and M. Mangano, Phys. Lett. 135B (1984) 445.

[18] J. Ellis, S. Ferrara and D.V. Nanopoulos, Phys. Lett. 114B (1982) 231;
W. Buchmüller and D. Wyler, Phys. Lett. 121B (1983) 321;
J. Polchinski and M.B. Wise, Phys. Lett. 125B (1983) 393.

[19] M.J. Duncan, Nucl. Phys. B221 (1983) 285;
J.F. Donoghue, H.P. Nilles and D. Wyler, Phys. Lett. 128B (1983) 55.

[20] J.-M. Gérard, W. Grimus, A. Masiero, D.V. Nanopoulos and Amitava Raychaudhuri, CERN preprint TH.3920.

[21] S.-P. Chia, Phys. Lett. 130B (1983) 315.

[22] Y. Meurice, Ann. Soc. Scient. Bruxelles 96 (1982) 28.

[23] M.A. Shifman, A.I. Vainshtein and V.I. Zakharov, Pis'ma Zh. Eksp. Teor. Fiz. 22 (1975) 123 [JETP Lett. 22 (1975) 55]; Zh.

Eksp. Teor. Fiz. 72 (1977) 1275 [Sov. Phys. JETP 45 (1977) 670]; Nucl. Phys. B120 (1977) 315.

[24] J.F. Donoghue, Univ. of Massachusetts preprint UMHEP-196 (1984).

[25] A.I. Sanda, Phys. Rev. D23 (1981) 2647;
N.G. Deshpande, Phys. Rev. D23 (1981) 2654.

[26] P. Langacker and B. Sathiapalan, Univ. of Pennsylvania preprint UPR-0251T.

[27] F.J. Gilman and J.S. Hagelin, Phys. Lett. 126B (1983) 111.

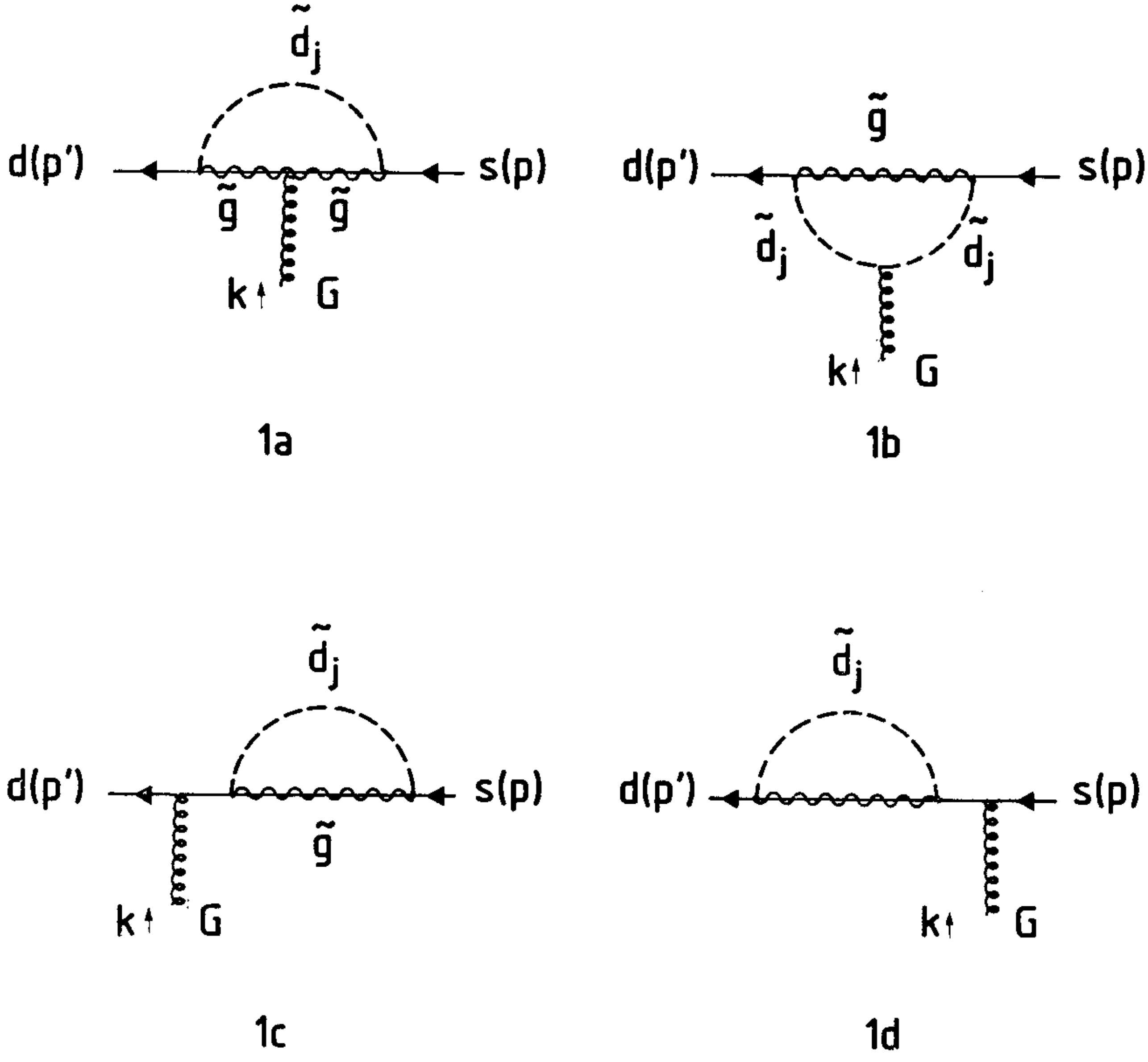
Figure Captions

1. The pengluino graphs.

2. Feynman rules for the vertices appearing in Fig. 1.

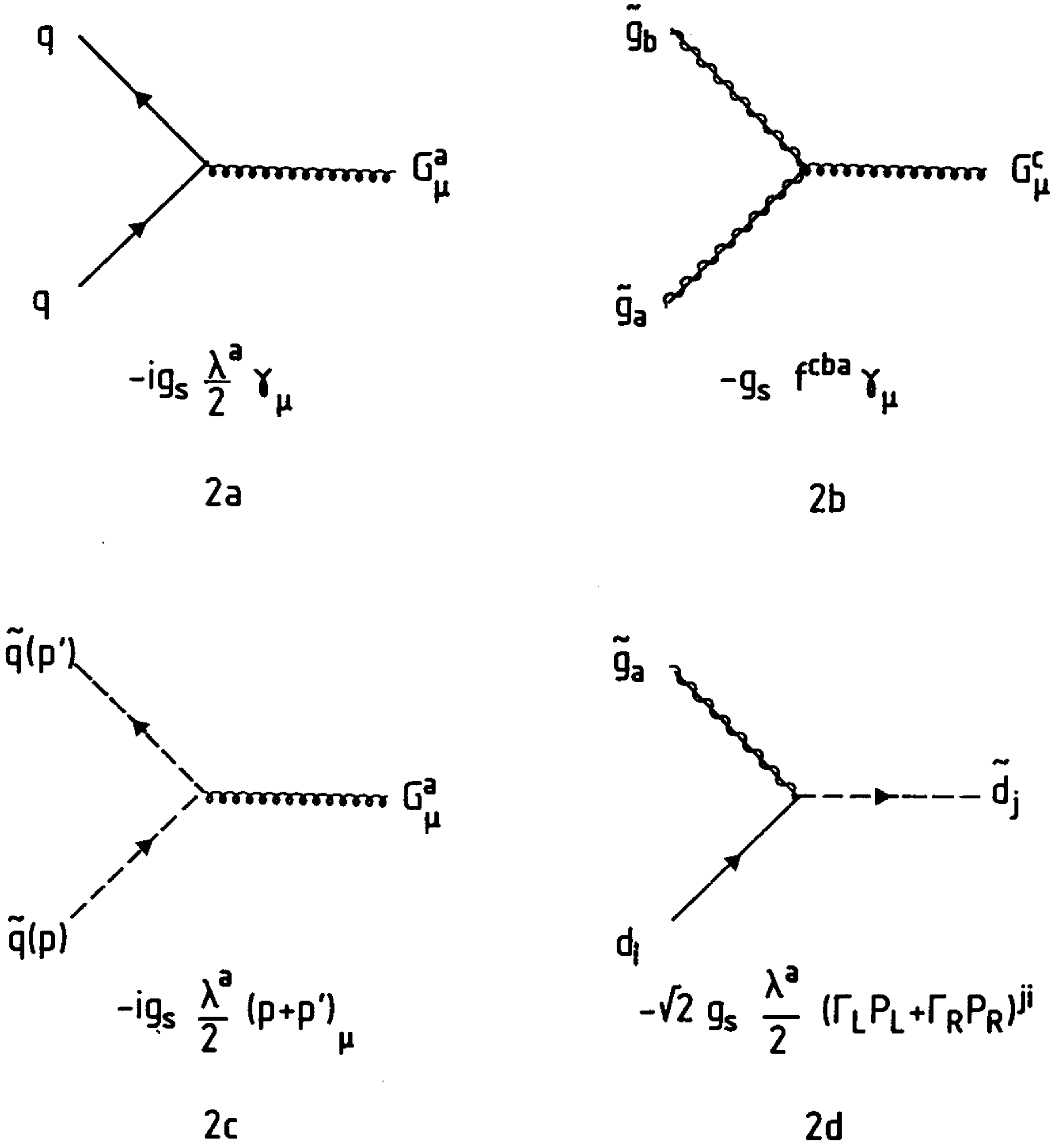
3. ε'/ε as a function of the top quark mass in the standard model for β=0.5 (dotted lines) and in the SUSY model (characterized in the text) for β=0.5 (dashed-dotted lines), 1.0 (full lines) and 1.5 (dashed lines). The bottom lifetime $\tau_B=\beta\times10^{-12}$sec. In this plot we choose the flavour violation parameter c=-1 and the $\tilde{d}_L$-$\tilde{d}_R$ mixing parameter $|A|$=3.

4. ε'/ε as a function of the bottom lifetime τ_B for a top quark mass equal to 40 GeV in the SUSY model (for different choices of the parameters c and A) and the standard model.



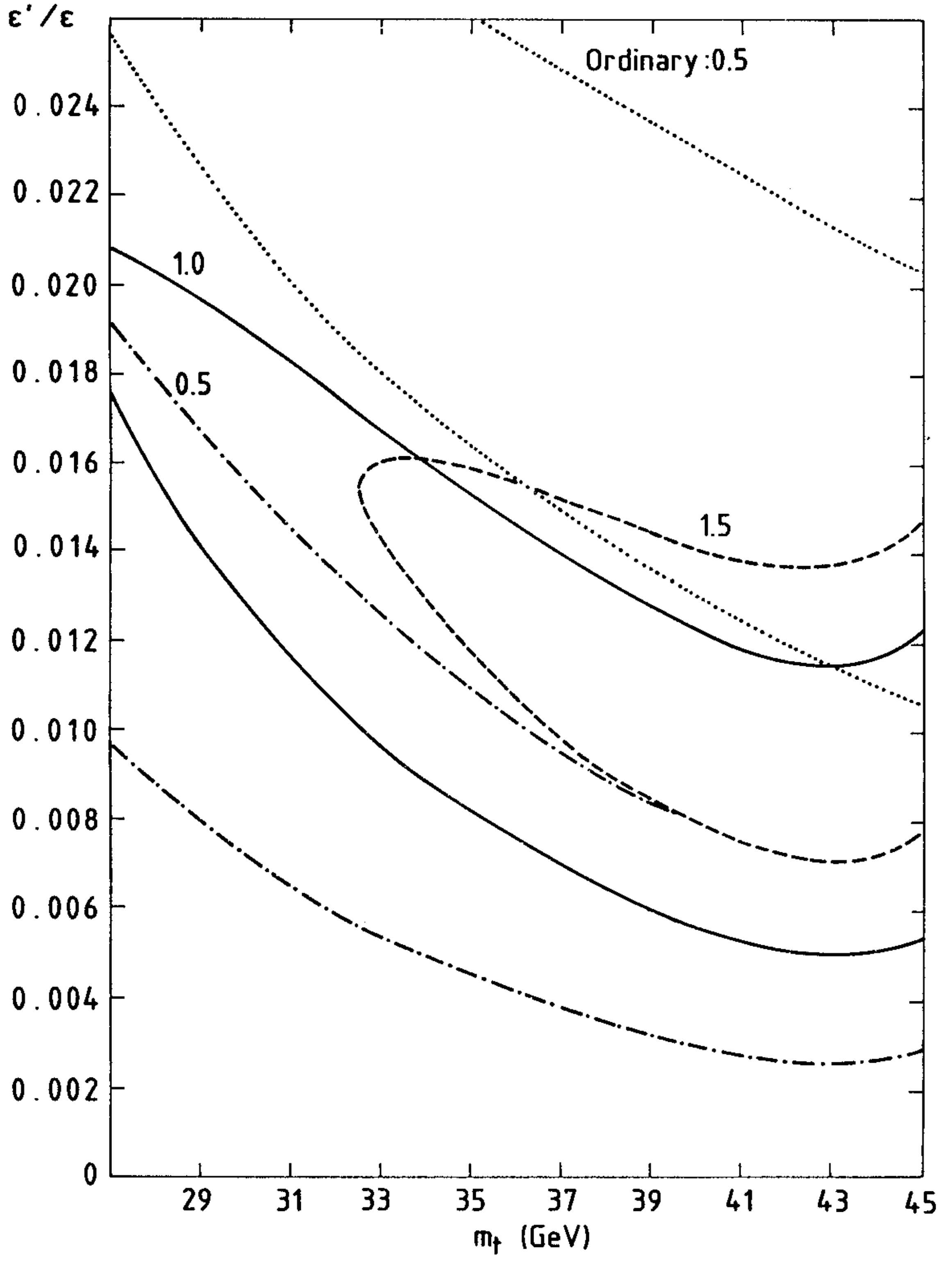
$\tilde{d}_j$
d(p')
s(p)
$\tilde{g}$
$\tilde{g}$
k
G
1a
$\tilde{g}$
d(p')
s(p)
$\tilde{d}_j$
$\tilde{d}_j$
k
G
1b
$\tilde{d}_j$
d(p')
s(p)
$\tilde{g}$
k
G
1c
$\tilde{d}_j$
d(p')
s(p)
k
G
1d

Fig. 1



q
q
G_μ^a
$-ig_s \frac{\lambda^a}{2} \gamma_\mu$
2a
$\tilde{g}_b$
$\tilde{g}_a$
G_μ^c
$-g_s \, f^{cba} \gamma_\mu$
2b
$\tilde{q}(p')$
$\tilde{q}(p)$
G_μ^a
$-ig_s \frac{\lambda^a}{2} (p+p')_\mu$
2c
$\tilde{g}_a$
$\tilde{d}_j$
d_i
$-\sqrt{2}\, g_s \frac{\lambda^a}{2} (\Gamma_L P_L + \Gamma_R P_R)^{ji}$
2d

Fig. 2



ε' / ε
Ordinary : 0.5
1.0
0.5
1.5
0.024
0.022
0.020
0.018
0.016
0.014
0.012
0.010
0.008
0.006
0.004
0.002
0
29
31
33
35
37
39
41
43
45
m_t (GeV)

Fig. 3

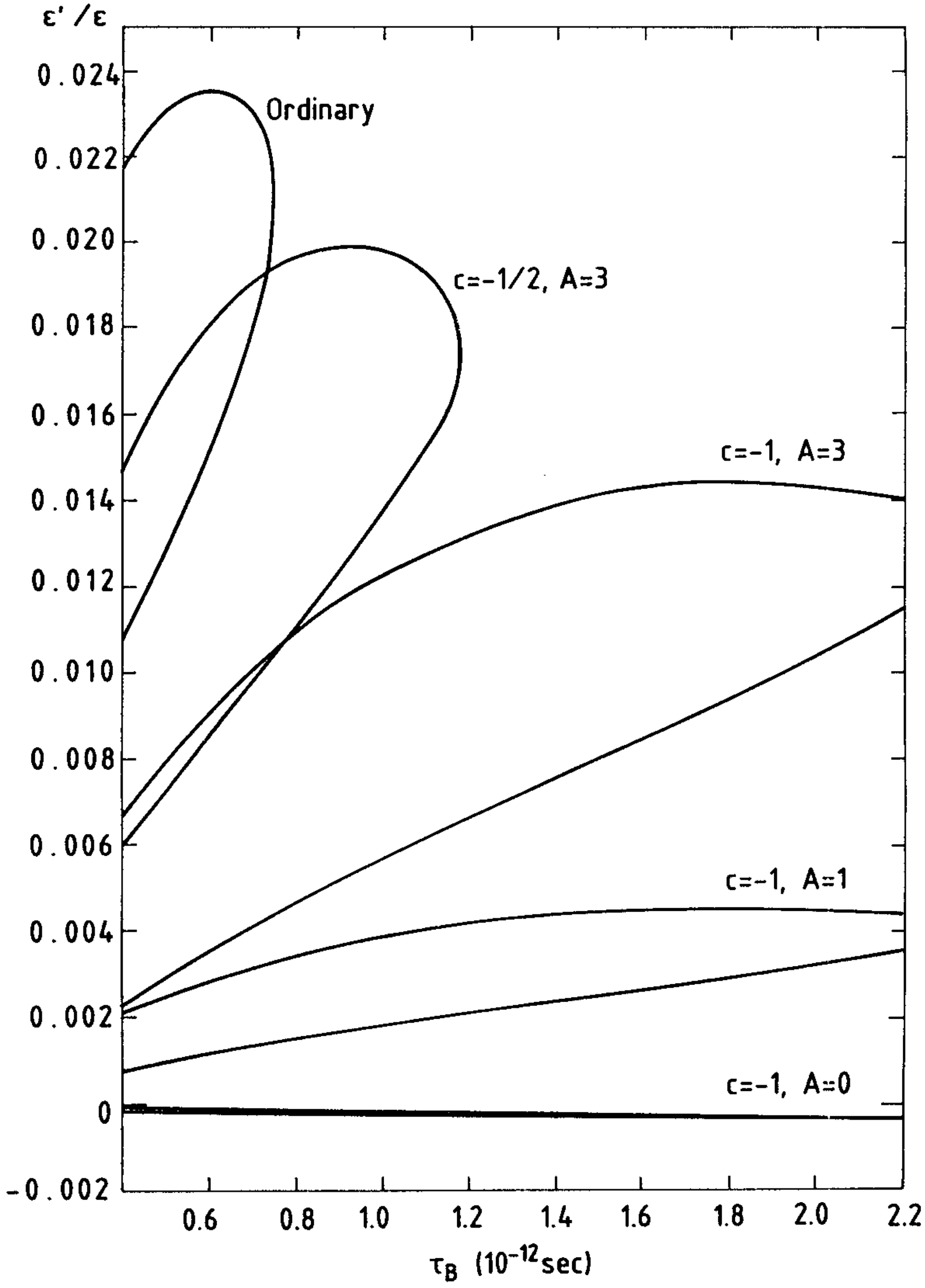

ε' /ε
0.024
0.022
0.020
0.018
0.016
0.014
0.012
0.010
0.008
0.006
0.004
0.002
0
-0.002
Ordinary
c=-1/2, A=3
c=-1, A=3
c=-1, A=1
c=-1, A=0
0.6
0.8
1.0
1.2
1.4
1.6
1.8
2.0
2.2
τB (10-12sec)

Fig. 4