$B^0-\bar{B}^0$ MIXING WITHIN AND BEYOND THE STANDARD MODEL

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ABSTRACT

The standard model with three fermion families is found to be compatible with the Argus observation of a relatively large amount of $B_d^0-\bar{B}_d^0$ mixing provided that the top quark is rather heavy. From an analysis of the existing information on mixing angles and of the constraints imposed by the kaon system, in particular by the ϵ parameter, we conclude that almost certainly $m_t \gtrsim 45$ GeV and probably $m_t \gtrsim 90$ GeV. In view of the Argus result the standard model unambiguously leads to the prediction of a nearly maximal amount of $B_s^0-\bar{B}_s^0$ mixing. Apart from the rather obvious case of a fourth family of quarks, most "minimal" extensions of the standard model preserve the prediction of a large $B_s^0-\bar{B}_s^0$ mixing. We discuss $B^0-\bar{B}^0$ mixing in minimal supersymmetric models compatible with UA1 lower bounds on gluino and squark masses. Contributions from charged Higgs exchange in the box diagrams are also discussed. While supersymmetry and charged Higgses can lead to an appreciable enhancement of $B^0-\bar{B}^0$ mixing, minimal left-right symmetric models actually predict a modest suppression of the effect with respect to the standard model.

1. Introduction

Recently, a nonvanishing amount of B^0_d - $\bar{B}^0_d$ mixing has been measured at DORIS by the Argus collaboration [1] (where B^0_d (B^0_s) is the neutral pseudoscalar meson made up of $b\bar{d}$ ($b\bar{s}$) quark-antiquark pairs). Evidence for non-vanishing B^0 - $\bar{B}^0$ mixing was previously also found by the UA1 collaboration [2] at CERN. Defining as usual [3] $r = P(B^0 \rightarrow \bar{B}^0)/P(B^0 \rightarrow B^0)$ ---which coincides with $\bar{r} = P(\bar{B}^0 \rightarrow B^0)/P(\bar{B}^0 \rightarrow \bar{B}^0)$ when the presumably small effect of CP violation is neglected---the Argus group finds consistent signals of mixing which lead to a combined result

$$r_d \approx 0.21 \pm 0.08 . \quad (1)$$

In particular, $r_d > 0.09$ at the 90% confidence level. This result is somewhat surprising as the central value of r_d in the standard model with three generations and a light top quark is much smaller [4]. The simplest interpretation of the Argus result in this context is that the top quark is rather heavy. A bound on the top quark mass of about $m_t \gtrsim 50$ GeV was indicated by the Argus collaboration and confirmed in ref. [5]. In the present paper we study the implications of the Argus result in the minimal standard model and beyond.

We start, in the framework of the minimal standard model, by summarizing the available information on the values of the relevant parameters. In particular, we study in detail the implications of the measured value of ϵ , the CP violation parameter of the kaon system [3], on the predicted value of r_d . In fact the choice of the CP violating phase that makes r_d maximum for a given m_t would correspond to $\epsilon=0$, so that a compromise is necessary.

After stating our conclusions on m_t and the resulting predictions on r_s in the standard model with three fermion families, we enlarge the discussion to include effects from possible new physics. We are interested in studying those mechanisms that could lead to an increase of r_d with respect to the prediction of the minimal standard model so that the observed value of r_d could be made compatible with a larger interval of values for the top quark mass. We briefly consider the obvious possibility of a fourth fermion family. We then analyze in detail the case of supersymmetric models, with soft breaking induced by gravity [6], which have been extensively studied in the past few years [7]. It has been pointed out [8] that in these models large contributions to $B^0-\bar{B}^0$ mixing are possible. However, the experimental lower limits on the masses of the relevant supersymmetric particles (gluinos and squarks) have recently been considerably improved at the CERN Sp $\bar{p}$ S collider by the UA1 collaboration [9]. We thus discuss the supersymmetric contributions to $B^0-\bar{B}^0$ mixing compatible with $K^0-\bar{K}^0$ physics, and with the present bounds on supersymmetric particles. Then, as in all supersymmetric models physical charged Higgs bosons are necessarily predicted [7], we also consider the possible effects (on $B^0-\bar{B}^0$ mixing) of box diagrams involving charged Higgs exchange [10,11]. This analysis is of

course interesting in itself independently of supersymmetry. We conclude that supersymmetry and charged Higgses can in fact lead to a sizeable increase of the predicted values of r_d provided that the gluinos, squarks, and/or charged Higgses are sufficiently light, and that the relevant couplings are in a specified domain of values. Finally we consider minimal left-right symmetric models [12] and conclude that these possibilities [13] do not appear promising in the present case. For example, the contribution of right handed W's in the box diagrams not only is small if the restrictive constraints from $K^0-\bar{K}^0$ physics are imposed [14], but partially cancels the standard model contribution.

2. Minimal Standard Model

In this section we shall set forth the basic formulae and discuss $B^0-\bar{B}^0$ mixing in the standard model with three generations. For B^0 mesons the mixing parameter can be written in the form [3]

$$r = \frac{P(B^0 \rightarrow \bar{B}^0)}{P(B^0 \rightarrow B^0)} = \frac{x^2 + y^2}{2 + x^2 - y^2} \simeq \frac{x^2}{2 + x^2} \quad (2)$$

where $x = \Delta M/\Gamma$, $y = \Delta\Gamma/2\Gamma$ with $\Delta M = M_L - M_S$, $\Delta\Gamma = \Gamma_L - \Gamma_S$ and $\Gamma = (\Gamma_S + \Gamma_L)/2$ (L and S refer as usual to the long and short lifetime mass eigenstates respectively). In eq.2 y^2 can be safely neglected for B mesons. For heavy quarks the short distance approach is considered to be particu-

larly reliable. In the standard model with three generations, $x_{d,s}$ are dominated by the top quark contribution to the box diagrams. Thus, one obtains [15]:

$$X_d = \frac{G_F^2}{6\pi^2} m_t^2 \tau_{B_d} B_{B_d} f_{B_d}^2 m_{B_d} |V_{td} V_{tb}^*|^2 \frac{A(z_t)}{z_t} \eta_t \quad (3)$$

where q stands for d or s . In the above equation G_F is the Fermi constant and $\tau_B = 1/\Gamma_B$ the B^0 lifetime. We use [9] $\langle B_d^0 | J_\mu^\dagger J^\mu | \bar{B}_d^0 \rangle = \frac{4}{3} B_{B_d} f_{B_d}^2 m_{B_d}$ where J_μ is the relevant charged weak current for quarks, $m_{B_d} = 5.275$ GeV is the mass of B_d^0 , and $B_{B_d} = 1$ corresponds to the dominance of the vacuum intermediate state. $V_{q_a q_b}$ denote the Cabibbo-Kobayashi-Maskawa (CKM) matrix elements [16], and η_t is the QCD correction factor [15], which we take as $\eta_t \approx 0.85$. Defining $z_t = m_t^2/M_W^2$, the function [17]

$$\frac{A(z)}{z} = \frac{1}{4} + \frac{9}{4(1-z)} - \frac{3}{2} \frac{1}{(1-z)^2} - \frac{3}{2} \frac{z^2 \ln z}{(1-z)^3} \quad (4)$$

is a monotonic slowly decreasing function of m_t varying between 1 at $m_t=0$ down to 0.75 for $m_t=M_W$ and to 0.52 for $m_t=200$ GeV.

The value of $B_B^{1/2} f_B$ is one of the main sources of uncertainty in the present problem (we leave aside for the moment the difference between B_d^0 and B_s^0). One would expect that the vacuum insertion approximation becomes more accurate with the increase of the meson mass, so that B_B

should be closer to unity than B_K . f_B is best evaluated from QCD sum rules. One method followed by different groups [18] consistently leads to $f_B \approx 100$ -130 MeV. Another, in principle equivalent, method [19] appears to lead to larger values ($f_B \approx 190 \pm 30$ MeV). On the experimental side, from a search for $D^\pm \rightarrow \mu^\pm \nu$ decays the MarkIII collaboration at SLAC obtains [20] the bound $f_D < 340$ MeV for charmed mesons. If scaled down with the asymptotic (for heavy quarks) law $f_B \sim m_B^{-1/2}$ the bound on f_D is translated into $f_B < 200$ MeV. This scaling law is valid for a non-relativistic picture of heavy quark-light quark mesons [21] (which is presumably adequate for f_D and f_B , but not necessarily for f_K) and is also supported by QCD sum rules [18,19]. Combining the available theoretical and experimental information we choose in the following

$$B_B^{1/2} f_B = (140 \pm 40) \text{ MeV} . \quad (5)$$

This is meant to imply that the whole range of conceivable possibilities is included within twice the quoted "standard" error. For $V_{q_a q_b}$ we take the approximate parametrization [22]

$$V = \begin{pmatrix} V_{ud} & V_{us} & V_{ub} \\ V_{cd} & V_{cs} & V_{cb} \\ V_{td} & V_{ts} & V_{tb} \end{pmatrix} \simeq \begin{pmatrix} 1 - \frac{1}{2}\lambda^2 & \lambda & A\lambda^3 \rho e^{i\varphi} \\ -\lambda & 1 - \frac{1}{2}\lambda^2 & A\lambda^2 \\ A\lambda^3 (1 - \rho e^{-i\varphi}) & -A\lambda^2 & 1 \end{pmatrix} \quad (6)$$

This expression is valid when terms proportional to λ^4 are neglected, which is completely adequate for our purposes. The parameter λ is to excellent approximation equal to the sine of the Cabibbo angle θ_c , or $\lambda = 0.221$. The experimental constraints [23] on the remaining parameters A , ρ and ϕ

are crucial for the present analysis, and are discussed at some length in what follows.

The present constraints on ρ are obtained from the experimental limit on $R = \Gamma(b \rightarrow u) / \Gamma(b \rightarrow c)$. The derivation of a limit on R from the actual data necessarily involves some model dependence. An accurate treatment of QCD corrections and phase space effects, together with a more model dependent but relatively secondary description of nonperturbative effects in semileptonic charm and beauty decays, was given in ref. [24]. When precision studies of semileptonic B-meson decays were first done, by the experiments CUSB and CLEO at CESR [25], the results of the above analysis were used to obtain from the data stringent bounds on R . Later, the bounds on R were considerably relaxed [26], on the basis of new data and a rather pessimistic attitude toward the degree of model dependence contained in the calculations of ref. [24]. New calculations were done [27]; as summarized in ref. [28] these results are in reasonable agreement among themselves and with ref. [24]. All recent theoretical analysis consistently indicates that indeed $R \lesssim 0.04$. However, for completeness in the following we present results for the conservative limit

$$R \equiv \Gamma(b \rightarrow u) / \Gamma(b \rightarrow c) < 0.08 \quad . \quad (7)$$

together with those obtained for smaller values of R . From the relation [24]:

$$(0.47 \pm 0.02) \cdot R = \left| \frac{V_{bu}}{V_{bc}} \right|^2 \quad (8)$$

we obtain

$$\lambda_f = \left| \frac{V_{bu}}{V_{bc}} \right| < 0.20 \quad \left(< 0.14 \right) \quad (9)$$

for $R < 0.08$ (< 0.04), so that

$$\rho \lesssim 0.9 \quad (\lesssim 0.6) \quad (10)$$

where ρ is defined in eq. 6.

The value of A is fixed by the B lifetime τ_B and semileptonic branching ratio B_{SL} through the relation

$$\tau_B |V_{bc}|^2 = \frac{B_{SL}}{Z_c} \frac{10^{-13} s}{(1 + R)} \quad (11)$$

A careful estimation of the parameter Z_c (proportional to $\Gamma(b \rightarrow c e \nu)$) was done in ref. [24], where the value $Z_c = 4.0 \pm 0.6$ was determined. Note that this formula is often used by inserting for Z_c a value obtained from the μ lifetime by simply scaling by the factor $(m_b/m_\mu)^5$, where some particular value (with no error) is inserted for m_b . This procedure obviously by far underestimates the error on $|V_{bc}|$. By inserting in eq. 11 $Z_c = 4.0 \pm 0.6$ and $B_{SL} = 0.117 \pm 0.006$ (taken from ref. [29] by averaging over the electron and muon branching ratios) one obtains for $R < 0.08$

$$\tau_B |V_{bc}|^2 = (2.9 \pm 0.6) \times 10^{-15} \text{ s} \quad (12)$$

Note that $x_{d,s}$ depend only on the combination $\tau_B A^2$ or $\tau_B |V_{bc}|^2$. Thus we do not need to specify τ_B and to take into account the corresponding error in the calculation of $B^0-\bar{B}^0$ mixing. However, for different purposes (e.g. the calculation of $|\epsilon|$), we use $\tau_B=1.11\pm0.16$ ps [30] and obtain

$$A = 1.05 \pm 0.17 \quad (13)$$

Since $V_{cb} \approx -V_{ts}$, the quantity $\tau_B |V_{ts}|^2$ is directly given in eq. 12; to obtain $\tau_B |V_{td}|^2$ we need to introduce $\cos \phi$:

$$\left| \frac{V_{td}}{V_{ts}} \right| = \lambda \sqrt{1 + \rho^2 - 2\rho \cos \phi} \quad (14)$$

The maximum value of this ratio (and hence of x_d and r_d) is obtained for $\cos \phi = -1$:

$$\left| \frac{V_{td}}{V_{ts}} \right| \lesssim \lambda (1 + \rho) \approx \frac{1}{2.4} \quad (15)$$

However, if $\cos \phi = -1$, then all CP violating effects would disappear. In particular, the parameter ϵ for kaons is linear in $\sin \phi$ [15]. Thus the measured value $|\epsilon| = (2.28 \pm 0.05) \times 10^{-3}$ imposes a limit on how closely $\cos \phi$ can approach -1. In order to implement this constraint we start from the well-known expression for $|\epsilon|$ [15]:

$$\begin{aligned}
|\epsilon| &= \frac{\text{Im } M_{12} + \frac{\xi}{2} (2 \text{Re } M_{12})}{\sqrt{2} \Delta m_K} = \\
&= \frac{G_F^2 M_W^2}{12 \sqrt{2} \pi^2} \frac{B_K f_K^2 m_K}{\Delta m_K} (D + 2 \xi F)
\end{aligned} \tag{16}$$

The kernels D and F are specified in eq.s 2.12 and 2.13 of the first of ref. [15]. D is proportional to $\text{Im}(V_{td}^* V_{ts}) = A^2 \lambda^5 r \sin \phi$. ξ is also proportional to $\sin \phi$. We take, following ref. [15],

$$\xi = -.54 A^2 \lambda^4 \rho \sin \phi \quad . \tag{17}$$

The ξ term is relatively small, so that the uncertainties in its value are not significant. We will use $f_K = 160$ MeV, $\Delta m_K / m_K = 0.71 \times 10^{-14}$ and $B_K = 1$. The charm quark mass, which enters in D and F , is taken as $m_c = 1.4$ GeV. In computing CP violating effects the approximate expression for the CKM matrix could be improved to include some terms of order up to λ^5 [22]. These corrections were included in our numerical calculations, although they are insignificant.

The resulting bounds on $\cos \phi$ as functions of m_t for $B_K < 1$ and different values of A and ρ are shown in figs. 1a,b. The well known difficulty of reproducing the experimental value of $|\epsilon|$ with m_t small is apparent for small values of ρ (fig. 2) and becomes severe for $\rho < 0.5$, while for $\rho \approx 0.9$ the difficulty at most arises for A smaller than its maximum allowed value. In fact, if A and ρ are as large as possible within their present uncertainties, then at $m_t \approx 40$ GeV, $(\cos \phi)_{\min} \approx -0.94$, so that the corresponding constraint on V_{td} does not alter eq. 14 significantly for $m_t \gtrsim 40$ GeV. Note that the boundary curves for $\cos \phi$ are not symmetrical for $\cos \phi \rightarrow -\cos \phi$ because of the small dependence of $|\epsilon|$ on $\text{Re}(V_{td}^* V_{ts})$ which is also

contained in D (eq. 16). In fact, for $m_t \approx 40$ GeV, $(\cos \phi)_{\max} \approx 0.7$. In conclusion, by combining eq. 12 with $|V_{bc}| \approx |V_{ts}|$, eq. 15, and the above results on $\cos \phi$, we obtain for $m_t \approx 40$ GeV:

$$\tau_B |V_{td}|^2 = (3.3 \pm 2.7) \times 10^{-16} \text{ s} \quad . \quad (18a)$$

For m_t large, where the constraints from ε can be neglected, one instead obtains

$$\tau_B |V_{td}|^2 = (0.02 \text{ to } 6.2) \times 10^{-16} \text{ s} \quad . \quad (18b)$$

We have also considered the present experimental limits on ε'/ε and the corresponding bounds implied for ξ . For current values [29] of the experimental upper limits on $|\varepsilon'/\varepsilon|$ no significant restriction is obtained, taking into account the considerable theoretical uncertainty that is also involved.

In a rather broad range of m_t near $m_t \approx 40$ GeV we can write eq. 2 in the simple approximate form:

$$x_d \approx 0.15 \left(\frac{\tau_B |V_{td}|^2}{3.3 \cdot 10^{-16} \text{ s}} \right) \left(\frac{B_B f_B^2}{(0.14 \text{ GeV})^2} \right) \left(\frac{m_t}{40 \text{ GeV}} \right)^2 \quad (19)$$

The Argus result, eq. 1, implies $x_d = 0.73 \pm 0.18$. Taking the central values for $\tau_B |V_{td}|^2$ and $B_B f_B^2$ one would need $m_t \gtrsim 90$ GeV in order to obtain $x_d > 0.55$ (the variation of $A(z_t)/z_t$ given in eq. 4 was also taken into account). On the other hand, if $\tau_B |V_{td}|^2$ and $B_B f_B^2$ are both increased to $6 \times 10^{-16} \text{ s}$ and $(180 \text{ MeV})^2$ respectively then one needs $m_t \gtrsim 45$

GeV. The predicted values of x_d (obtained from the exact form given in eq. 3) as a function of m_t , for different choices of the relevant parameters, are shown in fig. 3, together with the Argus result.

If the bound on R could be improved with respect to the value $R < 0.08$ used in fig. 3, then the restrictions on m_t in order to reproduce the observed value of x_d would become more severe. This is shown in fig. 4, where smaller values of R are considered for the same (extremal) values of the remaining parameters. In both figs. 3 and 4 we have included the effect of the bounds on $\cos\phi$ coming from $|\varepsilon|$.

While the absolute values of x_d and x_s are subject to considerable uncertainty, their ratio is instead fixed to within $SU(3)_{\text{flavour}}$ breaking corrections to be:

$$\frac{x_d}{x_s} = \left| \frac{V_{td}}{V_{ts}} \right|^2 (1 + o(SU(3) \text{ breaking})) \quad (20)$$

We also note that the $SU(3)$ breaking correction is presumably negative because one expects $f_{B_d} < f_{B_s}$ and $\tau_{B_d} < \tau_{B_s}$. Neglecting the $SU(3)$ breaking correction in eq. 20, and using the limit in eq. 15 one obtains:

$$\left| \frac{V_{td}}{V_{ts}} \right|^4 < \left(\frac{1}{2.4} \right)^4 = \frac{1}{33} = \frac{1}{a} \quad (21)$$

$$r_s > \frac{a r_d}{a r_d + 1} \quad (22)$$

so that $r_d > 0.13$ implies $r_s > 0.81$. The constraint in eq. 22 is shown in fig. 5. It is clearly important to measure r_s and check whether it is indeed close to one as required by eq. 22. There are measurements by the MARKII collaboration [31] that could pose a potential problem for the standard model with three generations. However, the MARKII limit is only dangerous if combined with questionable assumptions on the probability for the b quark to combine with a d or s antiquark to make a B^0_d or B^0_s meson.

The prediction for r_s given r_d is also true in many simple extensions of the standard model where the CKM matrix suppression is maintained. This is the case for the models we will consider in the latter part of this paper, based on supersymmetry, multiple Higgs doublets, and minimal left-right symmetry.

The simplest explanation of the violation of the inequality in eq. 22 would be the existence of a fourth family. In the presence of a fourth generation, because of the more relaxed unitarity constraints on the CKM matrix elements, even the top quark contribution by itself (i.e. without invoking the additional terms from the new up-type quark t') can lead to drastically different values of both the ratio x_d/x_s and the quantity x_d . For example, the analysis of mixing angles in ref. [23] leads to $V_{td} \lesssim 0.017$ (note that eq.18b corresponds to $V_{td} \lesssim 0.026$) for three generations. For four generations, their bound on V_{td} is relaxed by almost one order of magnitude-- $V_{td} \lesssim 0.15$ --which implies that practically all values of r_d are possible.

We conclude that the minimal standard model with three generations is compatible with the Argus result provided that the top quark mass is rather heavy, i.e., certainly $m_t \gtrsim 45$ GeV and probably $m_t \gtrsim 90$ GeV. But then r_s must be close to unity. For $r_d > 0.13$ one obtains $r_s \gtrsim 0.81$. It is very important to check this prediction. Moreover, the importance of measuring $R \equiv \Gamma(b \rightarrow u)/\Gamma(b \rightarrow c)$ is once more confirmed, as the constraints on the standard model are dramatically more severe for smaller R .

3. Supersymmetric models

In this section we discuss $B^0-\bar{B}^0$ mixing in minimal models where supersymmetry is broken softly by gravity [6,7], following the previous works in refs. [8]. The dominant additional contribution to $B^0-\bar{B}^0$ mixing from virtual exchange of supersymmetric particles arises from box diagrams with gluino and down squark exchange, as shown in fig. 6a. In fact, it is well known [32] that flavour changing couplings $\tilde{g} - q - \tilde{q}$ (where $\tilde{g}$ and $\tilde{q}$ denote gluinos and squarks) are induced by charged Higgsino exchange in (left-handed) down squark self mass diagrams (fig 6b). The corresponding down squark mass squared matrix acquires a component proportional to M_u^2 :

$$m_{\tilde{d}_L}^2 = \mu^2 \mathbb{1} + M_d^2 - |c| V^\dagger M_u^2 V \quad (23)$$

where $M_{d,u}$ are the down, up quark mass matrices. We are systematically neglecting nondiagonal terms in the $\tilde{d}$ mass squared matrix connecting left

and right squarks, since these presumably small effects can well be absorbed, at this exploratory stage, in our ignorance of the other parameters. The parameter $|c|$ plays a crucial role in the following in that the interesting effect vanishes for $c = 0$ and increases rapidly with $|c|$. c is related to the Yukawa couplings of the Higgs particles and is expected to be of $O(1)$. In various simple models [8] $|c|$ is found to be in the range 0.1 to 1. However, there are no stringent arguments that prevent somewhat larger values of $|c|$. Thus it appears reasonable to set $|c|=1$ if we are interested in obtaining an estimate of how large the supersymmetric contribution can be for a given value of m_t . Results for $|c| = 1/2$ will also be given.

The gluino box diagrams lead to the following expression [8] for ΔM_{ab} (QCD correction factors are not included at this stage):

$$\Delta M_{ab}^{\tilde{g}} = \frac{\alpha_s^2}{54 m_{\tilde{g}}^2} B_B f_B^2 m_B \sum_{ij} \lambda_{ab}^i \lambda_{ab}^j S_{ij}(z_i, z_j) \quad (24)$$

Here α_s is the QCD coupling at a scale of order $m_{\tilde{g}} \approx M_W$. We set in the following $\alpha_s \approx 0.12$. $\lambda_{ab}^i = V_{ia}^* V_{ib}$ and the index i can be u , c , or t . $S_{ij}(z_i, z_j)$ is a kernel given by

$$S_{ij} = 11 K_{ij} + 4 I_{ij} \quad (25)$$

with:

$$\begin{aligned} K_{ij} &= \frac{1}{z_i - z_j} \left[\frac{z_i^2 \ln z_i}{(1 - z_i)^2} + \frac{1}{1 - z_i} - (i \leftrightarrow j) \right] \\ I_{ij} &= \frac{1}{z_i - z_j} \left[\frac{z_i \ln z_i}{(1 - z_i)^2} + \frac{1}{1 - z_i} - (i \leftrightarrow j) \right] \end{aligned} \quad (26)$$

where $z_i = m_{\tilde{d}_i}^2 / m_{\tilde{g}}^2$ ($\tilde{d}_i = \tilde{d}, \tilde{s}, \tilde{b}$ are the down type squarks). In the excellent approximation (0.1%) of neglecting all quark masses except m_t one has $m_{\tilde{d}}^2 = m_{\tilde{s}}^2 = m_{\tilde{b}}^2 + |c|m_t^2$. Then, by using the unitarity of the CKM matrix, one can reduce eq. 24 to the simpler, more explicit form:

$$x_d^{\tilde{g}} = \frac{\alpha_s^2 \tau_{B_d} B_B f_B^2 m_{B_d}}{54 m_{\tilde{g}}^2} |V_{td} V_{tb}^*|^2 \Delta S_t(m_{\tilde{d}}, m_{\tilde{b}}, m_{\tilde{g}}) \quad (27)$$

where

$$\Delta S_t = S_{33}(z_3, z_3) + S_{11}(z_1, z_1) - 2S_{13}(z_1, z_3) \quad (28)$$

Other supersymmetric contributions arising from wino or Higgsino exchange are much suppressed, since α_s is relatively large and up squarks are nearly degenerate in mass.

In order to obtain numerical results, we need to input values for the gluino and squark masses. Taking into account the UA1 limits as presented in ref. [9] we consider $m_{\tilde{g}} \gtrsim 55 - 60$ GeV and $m_{\tilde{q}} \gtrsim 60 - 70$ GeV. For completeness we also include the results for the controversial, but nonetheless interesting, window for light gluinos, i.e., $m_{\tilde{g}} = 3$ to 5 GeV and $m_{\tilde{q}} > 100$ GeV. In figs. 7 and 8 we plot, for different values of the relevant parameters, the ratio $x^{\text{TOT}}/x^{\text{ST}}$ of the total resulting $x^{\text{TOT}} = x^{\text{ST}} + x^{\tilde{g}}$ for the $B_{d,s}^0$ mesons and the corresponding standard model contribution x^{ST} given in eq. 3. This ratio is particularly interesting because it is independent of the values of CKM angles, of $B_B f_B^2$ and of τ_B . As seen from figs. 7 and 8, the results for light or heavy gluinos are comparable for $m_t \sim 40$ GeV. Note that for small $m_{\tilde{g}}$ the explicit dependence on $m_{\tilde{g}}$ in

eq. 27 is cancelled by the asymptotic value of ΔS_t for large $z_{i,j}$. We see that the contribution of gluino box diagrams can indeed be important as r_d grows quadratically with x_d at small x_d . This is also apparent from figs. 9 and 10, where, as a function of m_t , the value of x_d including the supersymmetric effect is compared with the corresponding value in the standard model for different choices of the relevant parameters. It is to be stressed, however, that the supersymmetric effect drops very rapidly with $m_{\tilde{g}}$ (at large $m_{\tilde{g}}$) and $m_{\tilde{q}}$. The dependence on $|c|$ is also very sharp (the ratio $x^{\tilde{g}}/x^{ST}$ is approximately of the form $R=|c| f(|c|m_t^2)$). Note that the requirement of positivity of $m_{\tilde{b}}^2 = m_{\tilde{q}}^2 - |c|m_t^2$ implies an upper bound on m_t for given values of $m_{\tilde{q}}^2$ and $|c|$. The drop at large m_t of the supersymmetric contribution visible in figs. 7 and 9 is due to this upper bound.

In conclusion, supersymmetry could in principle explain the large value of r_d at moderate values of m_t , but only if gluino and squark masses are close to the present experimental bounds. On the other hand, we have checked that in all cases studied here the supersymmetric contribution to $M_{K_L} - M_{K_S}$ is completely negligible. Also, in the presence of supersymmetric effects, the constraints arising from the observed value of $|\varepsilon|$ are further relaxed.

It is interesting to recall that the same flavour changing mechanism that contributes to $B^0 - \bar{B}^0$ mixing can lead to sizeable branching ratios for decay rates of the types $b \rightarrow sy$ or $b \rightarrow sg$ which could be large enough to be observed, as discussed in refs. [33].

4. Charged Higgs bosons

In all the models based on supersymmetry, physical charged Higgs bosons are also predicted. Additional box diagrams contributions to $B^0-\bar{B}^0$ mixing obtained by replacing one or both of the W lines by virtual charged Higgses are therefore also present. We shall now discuss these terms briefly. Of course, charged Higgs contributions to $B^0-\bar{B}^0$ mixing can occur independently of supersymmetry, so that the following discussion is of more general validity. In refs. [10,11] the box diagrams with charged Higgs exchange were computed, and their implications for $B^0-\bar{B}^0$ mixing were discussed. In a model with two Higgs doublets, the relevant couplings of charged Higgses to quarks are given by

$$H \simeq \frac{g_2 H^+}{2\sqrt{2}M_W} \bar{u}_R \frac{v_d}{v_u} M_u V d_L + \dots \quad (29)$$

where u and d indicate up- and down-type quarks respectively, g_2 is the SU(2) weak gauge coupling, M_u is the diagonal up quark mass matrix, and V is the CKM matrix. v_d and v_u are the Higgs vacuum expectation values of the Higgs doublets giving mass to the down-type and up-type quarks respectively; they are proportional to the down and up quark masses via the corresponding Yukawa coupling constants. The ellipsis indicates terms proportional to the down quark mass matrix which can be neglected in comparison with the top quark contribution. The exact result for x_d can be obtained from ref. [10]. An approximate form (adequate for the scope of the present discussion) of the correction to the standard result, due to charged Higgs exchange, is given by [11]

$$x_d = x_d^{ST} \left[1 + \frac{1}{4} \left(\frac{v_d}{v_u} \right)^4 \left(\frac{m_t}{m_H} \right)^2 \right] \quad (30)$$

The limits of validity of this simple approximation obtained for $m_t \ll m_H$, and neglecting the box diagram with one Higgs and one W propagator, are discussed in ref. [11]. For $m_t \lesssim m_H$, the exact expression gives x_d larger, by about 10% or less. One sees, for example, that for $m_H \sim M_W$ and m_t as low as $m_t \sim 30$ to 40 GeV a ratio $v_d/v_u \sim 2$ to 4 is sufficient to yield a possible explanation for the observed value of r_d . Such values for the ratio of vacuum expectation values are perfectly admissible (see ref. [11] for current limits). However, we note that the perhaps appealing requirement of using the duplication of doublets to decrease the differences among the Yukawa couplings of Higgs bosons to quarks would lead to $v_d/v_u < 1$ since up quarks are on the average heavier than down quarks.

5. Minimal left-right symmetric models

In this section we review B^0 - $\bar{B}^0$ mixing in the context of minimal left-right symmetric models [12] as discussed in refs. [13,14] and add some observations. By minimal models we mean both a single multiplet of Higgs bosons coupled to quarks (transforming as $(1/2, 1/2, 0)$ of $SU(2)_L \otimes SU(2)_R \otimes U(1)$) and spontaneous CP violation with identical or complex conjugate CKM matrices for left-handed and right-handed quarks. It is well known that the contribution to the K_L - K_S mass difference arising from (right-handed) W_R exchange in the box diagrams are much larger than for the ordinary (left-handed) W_L of the same mass [14]. Requiring that the short-distance contribution to $M_{K_L} - M_{K_S}$ does not exceed the experimental

value leads in fact to a strong limit on the W_R mass: $M_{W_R} > 1.5$ to 2.5 TeV [13,14]. There are also additional contributions to $M_{K_L} - M_{K_S}$ arising from tree-level diagrams with flavour changing Higgs exchange, with masses M_H of order M_{W_R} . Thus a lower limit on M_H of several TeV can be derived [13,14].

In the approximation of neglecting the small mixing ξ_{LR} between left and right components in the charged weak boson mass matrix, the charged current terms in the Hamiltonian can be written

$$H \approx \frac{g_2}{\sqrt{2}} \left\{ W_L^\dagger \bar{u} V_L \gamma_\mu \left(\frac{1-\gamma_5}{2} \right) d + W_R^\dagger \bar{u} V_R \gamma_\mu \left(\frac{1+\gamma_5}{2} \right) d \right\} + h.c. \quad (31)$$

where V_L and V_R are the CKM matrices for left and right handed quarks respectively, which we assume to be the same in the following. The total resulting box diagram contributions to $x_{d,s}$, including the standard contribution x^{LL} from W_L exchange plus the W_R and unphysical scalar exchange terms, is given by:

$$X_{d,s}^{TOT} = X_{d,s}^{LL} \left[1 + \frac{3}{2} \frac{M_{W_L}^2}{M_{W_R}^2} \left(\frac{M_{B_{d,s}}^2}{(m_b + m_{d,s})^2} + \frac{1}{6} \right) \cdot \frac{1}{\eta_t} \cdot \frac{z_t}{A(z_t)} \cdot \left[\eta_{t1}^{LR} \left(\frac{4-z_t}{1-z_t} + \frac{4-2z_t+z_t^2}{(1-z_t)^2} \ln z_t \right) + \eta_{t2}^{LR} \ln \frac{M_{W_L}^2}{M_{W_R}^2} \right] \right] \quad (32)$$

Here η_t , z_t , and $A(z_t)/z_t$ are defined in eqs. 3 and 4, and $\eta_{t1}^{LR} \approx \eta_{t2}^{LR} \approx 1.8$ are QCD correction factors. From the above formula one can immediately derive the important results that

a) the amount of $B^0-\bar{B}^0$ mixing is not very much altered with respect to the standard model result, once the constraints on M_{W_R} imposed by the observed value of $M_{K_L} - M_{K_S}$ are taken into account.

b) the sign of the effect for $m_t \lesssim 250$ GeV goes in the direction of reducing the amount of mixing with respect to the standard model

c) the absolute value of the correction slowly decreases with m_t .

In fact, compared with the $K^0-\bar{K}^0$ system, the importance of the diagrams with left-right exchange is much decreased because $(m_B^2/(m_b+m_{d,s})^2) \sim 1$ is small in comparison with $(m_K^2/(m_s+m_d)^2)$. Moreover, the dominant contribution for the $K^0-\bar{K}^0$ system arises from charm exchange. Then $\ln z_t \rightarrow \ln m_c^2/m_W^2$ in eq. 31 and the corresponding term is quite large. As a consequence, for values of M_{W_R} large enough to be compatible with $m_{K_L}-m_{K_S}$, the corresponding correction for $B^0-\bar{B}^0$ mixing is small. As already mentioned it is negative for all interesting values of m_t , i.e., $m_t \lesssim 250$ GeV, so that it tends to decrease the expected amount of mixing with respect to the prediction of the standard model. These features are evident from fig. 11 where the ratio $x_d^{\text{TOT}}/x_d^{\text{LL}}$ is plotted as a function of m_t for $m_{W_R} = 1.5$ and 2.5 TeV.

Similarly, the Higgs exchange term is also negative, and smaller than for the $K^0-\bar{K}^0$ system. Neglecting QCD corrections, one obtains:

$$X_{d,s}^{\text{Higgs}} = -X_{d,s}^{\text{LL}} \frac{\sqrt{2} 6 \pi^2}{G_F m_H^2} \left(\frac{m_{B_{d,s}}^2}{(m_b + m_{d,s})^2} + \frac{1}{6} \right) \frac{z_t}{A(z_t)} \quad (33)$$

Since the additional logarithmic enhancement due to the relative smallness of the charm mass is not present in this case, the absolute value of the Higgs correction can be larger than that arising from the box diagrams.

6. Summary and Conclusions

We summarize in this section our main results. In the standard model with three families the Argus measurements of $B_d^0-\bar{B}_d^0$ mixing indicates that the top quark is rather heavy, i.e. $m_t \gtrsim 45$ GeV. However, in order to obtain such a low value for m_t , all other badly known parameters (i.e., $B_B^{1/2} f_B$, $\cos \phi$, $R = \Gamma(b \rightarrow u)/\Gamma(b \rightarrow c)$) have to be tuned simultaneously at their theoretical or experimental boundary values which is rather unlikely. Thus we conclude that probably $m_t \gtrsim 90$ GeV. In the standard model with three fermion families the Argus value for r_d immediately implies that the $B^0-\bar{B}^0$ mixing is nearly maximal, i.e. $r_s \gtrsim 0.81$. It is very important to check this prediction which is completely independent of the unknown values of m_t and $B_B^{1/2} f_B$, and only follows from the observed CKM matrix element structure, and its three-by-three unitary nature. If r_s is found to be smaller than expected, the most direct indication would be for a fourth fermion family.

On the other hand, many simple forms of new physics preserve the CKM matrix structure and therefore the prediction of a large value of r_s . We have studied in detail the case of minimal supersymmetric models and that of charged Higgs exchange. Given the present bounds on the masses of gluinos and squarks, the supersymmetric contributions can only be important (for a relatively light top quark) if gluinos and squarks are close to their minimum allowed values, and the parameter $|c|$ in eq. 23 is close to unity or larger. Similarly, the contribution of charged Higgses can be important for small M_H , or if $M_H \approx M_W$, provided that the vacuum expectation value proportional to down-type quark masses is larger than that pro-

portional to up-type quark masses---a perhaps counter intuitive but otherwise fully allowed possibility. Both the supersymmetric and the charged Higgs contribution rapidly increase with m_t at fixed values of the new particle masses and couplings. While supersymmetry and charged Higgses can lead to an enhancement of B^0 - $\bar{B}^0$ mixing, this is not the case for minimal left-right symmetric models (with equal or complex conjugate CKM matrices for left- and right-handed quarks). The effect of right-handed W 's and flavour-changing Higgses would tend to suppress by a small amount the predicted value of r_d and r_s with respect to the standard model.

Finally the importance of better measurements of R is once more emphasized. With better bounds on R , the minimum value of m_t compatible with the measured value of r_d increases rapidly. On the other hand, once m_t is large enough to accommodate r_d , then the constraints following from the experimental value of $|\varepsilon|$ for the kaon system are automatically satisfied.

While writing up this article, we received some related papers on B^0 - $\bar{B}^0$ mixing in the standard model [34].

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FIGURE CAPTIONS

1. Limits on $\cos\phi$ obtained from the experimental value of the CP violating parameter $|\epsilon|$ for the kaon system, as functions of the top quark mass, m_t , for various values of p . The solid (dashed) lines include (do not include) the effect of the ξ term, given in eqs. 16 and 17. In figure (a), we have taken $A = 1.05$ (the central value in eq. 13) and in (b), $A = 1.22$ (the maximum value). The parameters ϕ , p and A are defined in eq. 6. The indicated values of $p = 0.9$, 0.6 and 0.3 correspond to $R = \Gamma(b \rightarrow u) / \Gamma(b \rightarrow c) \approx 0.08$, 0.04 , and 0.009 respectively.
2. Lower limits on the top quark mass m_t from the observed value of the CP violating parameter $|\epsilon|$ for the kaon system as a function of p for $B_k = 2/3, 1$. The solid (dashed) lines include (do not include) the effect of the ξ term, given in eqs. 16 and 17. The curves are obtained for $A = 1.22$ (the maximum value allowed in eq. 13). The parameters A and p are defined in eq. 6; p is related to $R = \Gamma(b \rightarrow u) / \Gamma(b \rightarrow c)$ by eqs. 7, 8 and 9.
3. Predicted values of $B^0_d - \bar{B}^0_d$ mixing in the standard model with three fermion families, as a function of the top quark mass, for different choices of the relevant parameters. The Argus value (eq. 1) for r_d is also shown, which translates to $x_d = 0.73 \pm 0.18$ by the relation $r_d = x_d^2 / (2 + x_d^2)$. All curves are for $p < 0.9$ (or $\Gamma(b \rightarrow u) / \Gamma(b \rightarrow c) < 0.08$) where p is defined in eq. 6. In the curves (a)-(e) we use the following values of the parameters $F = B_B^{1/2} f_B$, $\cos\phi$ and $T = \tau_B |V_{bc}|^2 / 10^{-15} \text{s}$ (see eqs. 5, 6, and 12):

- a. $F = 180 \text{ MeV}$, $T = 3.5$, and $\cos\phi = (\cos\phi)_{\min}$ (the largest negative value of $\cos\phi$ allowed by the $|\varepsilon|$ constraint). This curve gives the lower bound $m_t > 45 \text{ GeV}$ quoted in the text.
- b. $F = 140 \text{ MeV}$, $T = 3.5$, and $\cos\phi = (\cos\phi)_{\min}$.
- c. $F = 140 \text{ MeV}$, $T = 2.9$, and $\cos\phi = 0$. This curve corresponds to about the central prediction of the standard model, and leads to the statement in the text that probably $m_t > 90 \text{ GeV}$.
- d. $F = 100 \text{ MeV}$, $T = 2.3$, and $\cos\phi = (\cos\phi)_{\max}$ (the largest positive value of $\cos\phi$ allowed by the $|\varepsilon|$ constraint).
- e. $F = 100 \text{ MeV}$, $T = 3.5$, and $\cos\phi = (\cos\phi)_{\max}$.

Note that (e) actually lies below (d), because of the relaxed constraint on $(\cos\phi)_{\max}$.

4. Maximum value of x_d compatible with the standard model with three fermion families as a function of the top quark mass, for different values of ρ . ρ is related to $R = \Gamma(b \rightarrow u) / \Gamma(b \rightarrow c)$ by eqs. 7, 8 and 9. ($\rho = 0.9, 0.6$ and 0.3 correspond to $R \approx 0.08, 0.04$, and 0.009). The $\rho = 0.9$ curve corresponds to curve (a) of fig. 3, i.e., $B_B^{1/2} f_B = 180 \text{ MeV}$, $\tau_B |V_{bc}|^2 = 3.5 \times 10^{-15} \text{ s}$, and $\cos\phi = (\cos\phi)_{\min}$. The other curves are also given for this parameter set, where $(\cos\phi)_{\min}$ is the appropriate minimum value allowed by the $|\varepsilon|$ constraints, which changes with ρ .
5. The solid curve shows the minimum value of r_s that corresponds to each value of r_d in the standard model with three fermion families. $r_{d,s}$ are the mixing parameters of B_d^0 and B_s^0 mesons respectively. The bound on r_s given r_d follows from eqs. 20, 21 and 22, and is independent of m_t , $B_B^{1/2} f_B$ and τ_B . The shaded area is the allowed region implied by the Argus result, eq. 1.
6. a. Box diagram with gluino and squark exchange.

- b. Charged Higgsino and top quark exchange contributing to flavour non-diagonal b to d (left-handed) squark mass matrix elements.

7. The ratio $x^{\text{TOT}}/x^{\text{ST}}$ for B^0_d or B^0_s mesons, where $x^{\text{TOT}} = x^{\text{ST}} + x^{\tilde{g}}$ is the sum of the standard term, eq. 3, plus the supersymmetric contribution, given in eq. 27. The solid lines correspond to $|c| = 1$, while the dashed line is computed for $|c| = 1/2$, where $|c|$ is defined in eq. 23. The gluino and down squark masses (in GeV) are indicated on the curves.
8. The ratio $x^{\text{TOT}}/x^{\text{ST}}$ (see fig. 7) for light gluinos. The gluino and down squark masses (in GeV) are indicated. The parameter $|c|$ is taken to be 1.
9. Comparison of the predictions for $B^0_d - \bar{B}^0_d$ mixing of the standard model (dashed) with those (solid) of the supersymmetric model with the indicated values of the gluino and down squark masses (in GeV). The Argus result is also shown. The standard model curve corresponds to the same choice of parameters as in curve (c) of fig. 3. It approximately represents the central prediction of the standard model, as a function of the top quark mass m_t . The supersymmetric curves (obtained for $|c| = 1$) are calculated for the same values of τ_B , $|V_{bc}|^2$, $B_B^{1/2} f_B$ and $\cos\phi = 0$ as for the standard model.
10. Same as fig. 9, but for light gluinos.
11. The ratio $x^{\text{TOT}}/x^{\text{LL}}$, given in eq. 32, for $B^0_{d,s}$ mesons, as a function of m_{top} , for different values of the mass of the right-handed charged weak boson W_R , chosen near the lower bound allowed by the observed mass difference of K_L and K_S . The curves were obtained for $[m_{B_{d,s}}^2 / (m_b + m_{d,s})^2 + 1/6] \approx 1$.

Fig. 1a

Fig. 1b

Fig. 2

Fig. 3

Fig. 4

Fig. 5

(a)

(b)

Fig. 6

Fig. 7

Fig. 8

Fig. 9

Fig. 10

Fig. 11