Université de Montréal

BRISURE CP/T VIA LES PRODUITS TRIPLES DANS LES DÉSINTÉGRATIONS DES HADRONS À SAVEUR DE BEAUTÉ

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Thèse présentée à la faculté des études supérieures
En vue de l'obtention du grade de DOCTORAT
en PHYSIQUE

Mai, 2003

Université de Montréal
Faculté des études supérieures

Cette thèse intitulée:

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RÉSUMÉ

Cette thèse est constituée de quatre articles qui ont tous été publiés. Le sujet principal est l'étude d'asymétries de brisure CP/T via les produits triples dans les désintégrations des hadrons contenant un quark b (mésons B ou hypérons Λ_b) et ce, dans le modèle standard de la physique des particules et au-delà.

Dans le premier article, on se place au niveau des quarks. On étudie, dans le modèle standard, les asymétries de brisure T dues aux produits triples (PT) dans le processus inclusif $b \to su\bar{u}$. Ensuite on cite les applications possibles au niveau hadronique. Seuls deux produits triples donnent des asymétries assez larges. On trouve une asymétrie d'environ 5% pour $\vec{p}_u \cdot (\vec{s}_u \times \vec{s}_{\bar{u}})$ ou $\vec{p}_{\bar{u}} \cdot (\vec{s}_u \times \vec{s}_{\bar{u}})$ et entre 1% et 3% pour $\vec{s}_b \cdot (\vec{p}_u \times \vec{p}_s)$. Pour les autres PT présents dans l'amplitude au carré, les asymétries sont négligeables. Ces résultats permettent de tester le MS et de voir, s'il y a lieu, la manifestation de la physique au-delà du MS (la "nouvelle physique"). En effet, le premier type d'asymétries peut être recherché dans les désintégrations d'un méson B en deux bosons vecteurs et le deuxième type, dans les désintégrations d'un hypéron Λ_b en un hypéron Λ avec deux pions. Les asymétries des PT contenant le spin du quark s sont négligeables. La mesure d'un tel signal serait la preuve de l'existence d'une nouvelle physique.

Dans le deuxième article, on calcule les asymétries de brisure T dues aux produits triples dans les désintégrations non charmées de Λ_b, dans le modèle standard et en utilisant la méthode de factorisation. On s'intéresse aux modes de désintégation de Λ_b en un baryon léger de spin 1/2 (F) avec un méson soit pseudoscalaire (P) soit vecteur (V). On trouve une grande asymétrie (18%) pour le mode $\Lambda_b \to pK^-$ correspondant au mode inclusif $b \to su\bar{u}$. Pour les autres modes du type $\Lambda_b \to FP$ l'asymétrie est au plus de l'ordre de 5%. Pour

les modes $\Lambda_b \to FV$, les asymétries sont très petites, valant au plus 1%. Ceci pousse à la recherche d'une nouvelle physique donnant lieu a de plus grands signaux dans ces modes. C'est ce qui nous a encouragé à écrire le troisième article.

Le troisième article de cette thèse concerne la recherche de signaux de brisure T via les produits triples dans la nouvelle physique. On a construit un lagrangien efficace pour les desintégrations non charmées de Λ_b, contenant tous les termes possibles de dimension 6. On a calculé la contribution de ces termes aux asymétries en utilisant la méthode de factorisation. Cela nous a permis de voir quels opérateurs (donc quels modèles de la nouvelle physique) donnent de grandes asymétries T. Ceci est intéressant surtout lorsque les prédictions du MS sont modifiées. En effet des asymétries négligeables dans le MS sont augmentées par les nouveaux opérateurs à jusqu'à 50%. Ce calcul est applicable à n'importe quel modèle de la nouvelle physique et permet de faire un tri parmis les opérateurs en mesurant les signaux de PT dans les désintégrations des Λ_b. On a appliqué nos résultats à deux modèles: la supersymétrie avec brisure de la parité R et "FCNC's" avec comme médiateur le Z' leptophobique.

Le quatrième article constitue un complément au sujet principal de cette thèse. Il concerne l'étude d'asymétries de polarisation et "avant-arrière" ("for ward-backward") dans le mode $b \to s\tau^+\tau^-$, avec les deux leptons polarisés. Cela permet de tester le modèle standard et de sonder la nouvelle physique. Il y a un total de 31 asymétries dans ce mode dont 9 sont de l'ordre de 10% et plus, donc mesurables. Cela permet la mesure des coefficients de Wilson et de la masse du quark b et de tester ainsi le modèle standard. On trouve que certaines asymétrie sont nulles dans le MS. Donc si par des mesures on les trouve non nulles ou encore si d'autres sont trouvées plus grandes que les valeurs prédites, ce sera un signal de présence d'une nouvelle physique. Un exemple intèressant est l'existence d'asymétries dues à des produits triples (ce sont de faux signaux de brisure CP car dûs seulement aux phases fortes) qui sont négligeables dans le MS (de l'ordre de 10^{-2}). Leur observation constituerait un signal évident de la nouvelle physique.

MOTS CLÉS

MODÈLE STANDARD

NOUVELLE PHYSIQUE

SPIN

POLARISATION

PARITÉ

CONJUGAISON DE CHARGE

RENVERSEMENT DU SENS DU TEMPS

ASYMÉTRIE AVANT-ARRIÈRE

LAGRANGIEN EFFICACE

MODÈLE DE FACTORISATION

SUMMARY

This thesis is comprised of four articles, all of which have been published. The main subject is the study, in and beyond the standard model of particle physics, of CP/T-violation via triple products in decays of hadrons containing a b quark (B mesons or Λ_b hyperons).

In the first article, we study, within the standard model (SM), the inclusive mode $b \to su\bar{u}$. We look for triple product asymmetries (TP's), which break T (time reversal symmetry). We then mention possible applications at the hadron level. Only two TP's give non-negligible asymmetries. We find an asymmetry of about 5% for $\vec{p}_u \cdot (\vec{s}_u \times \vec{s}_{\bar{u}})$ or $\vec{p}_{\bar{u}} \cdot (\vec{s}_u \times \vec{s}_{\bar{u}})$, and between 1% and 3% for $\vec{s}_b \cdot (\vec{p}_u \times \vec{p}_s)$. The asymmetries are negligible for all other triple products. These results allow one to test the SM and to search for evidence of physics beyond the SM ("new physics"). Indeed, the first class of asymmetries can be searched for in decays of a B meson into two vector bosons and the second class in decays of a Λ_b hyperon into a Λ and two pions. Asymmetries of TP's containing the s-quark spin are negligible. Measuring such a signal would be an evidence of new physics.

In the second article, we calculate, within the standard model, T-violating triple product asymmetries in charmless Λ_b decays using the factorization model. We focus on modes where the Λ_b decays into a light spin-1/2 baryon (F) with either a pseudoscalar (P) or a vector (V) boson. We find a large asymmetry (18%) for the decay $\Lambda_b \to pK^-$, whose inclusive mode is $b \to su\bar{u}$. The asymmetry for the other modes of the class $\Lambda_b \to FP$ is at most 5%. For the modes $\Lambda_b \to FV$, the asymmetries are tiny, not exceeding 1%. This made us think about new physics operators which could give larger signals in these modes. That is the subject of the third article.

The third article of this thesis consists of the search for T-violating signals via triple products within the context of new physics. We constructed an effective lagrangian for charmless Λ_b decays, with all possible dimension-6 terms. Using factorization, we calculated the contribution of these terms to the asymmetries. This allowed us to see which operator (hence which new-physics model) leads to large asymmetries. The calculations are most interesting when the SM predictions are modified. Indeed, negligible asymmetries in the SM can be enhanced by the new operators to up to 50%. The results of this paper can be applied to any new physics operators and allow one to determine which operators contribute to T-violating signals and which do not. We have applied our results to two specific models: supersymmetry with R-parity breaking and leptophobic Z'-mediated FCNC's.

The fourth article is a complement to the main subject of this thesis. It studies polarization and forward-backward asymmetries (which are not CP/T-violating) in the decay $b \to s\tau^+\tau^-$, with both leptons polarized. Of a total of 31 asymmetries, 9 are measurable, that is, they exceed 10%. This study allows one to exhaustively test the SM exhaustively and to probe new physics by extracting the Wilson coefficients and the b-quark mass. Certain asymmetries vanish in the standard model. Hence, if one of them is measured to be nonzero, or if others are measured to be bigger than their standard-model values, that would be a clear signal of new physics. An interesting example is the existence of triple product asymmetries (these are fake CP-violating signals due to strong phases) which are negligible in the SM ($\sim 10^{-2}$). Their observation would be a clear signal of new physics.

KEY WORDS

STANDARD MODEL

NEW PHYSICS

SYMMETRY

TRIPLE PRODUCT

POLARIZATION

CP/T-VIOLATION

DECAY

FORWARD-BACKWARD ASYMMETRY

EFFECTIVE LAGRANGIAN

FACTORIZATION MODEL

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(À part la première, elles font toutes partie du quatrième article)

LISTE DES SIGLES ET ABRÉVIATIONS

BSS	Brisure spontanée de la symétrie
C	"Charge conjugation" ou Conjugaison de charge
CKM	Cabibbo-Kobayashi-Maskawa
FCNC	"Flavour-changing neutral current"
L	"Left-handed"
LSP	"Lightest supersymmetric particle"
MS	Modèle standard
MSSM	"Minimal supersymmetric model"
NP	"New physics"
P	"Parity" ou parité
PT	Produit triple
QCD	"Quantum chromodynamics"
R	"Right-handed"
RAMBO	"Random momenta beautifully organized"
SM	"Standard model"
RESP.	Respectivement
T	"Time reversal" ou renversement du sens du temps
TP	"Triple product"

DÉDICACE

À Mes très chers papa et maman.
À Noureddine, prunelle de mes yeux.
À mon cher frère Khaled, que Dieu le guérisse.
À mes six autres frères et soeurs, D, N, H, H, H, et A.
À tous les autres membres de ma famille, grands et petits.
À Mes adorables, Hafsa, Belkiss et Joman; toutes nées durant mes études de doctorat.

REMERCIEMENTS

Un grand MERCI à mon directeur de recherche, Monsieur David London.

Merci à Monsieur Alakabha Datta.

Merci à Madame Nita Sinha et à Monsieur Rahul Sinha.

Merci à tous mes professeurs.

Merci aux administrateurs et aux techniciens du laboratoire René J.-A. Lévesque.

Merci aux administrateurs du département de physique.

Merci à tous ceux qui ont aidé, de près ou de loin, à l'accomplissement de ce travail.

AU NOM DE DIEU LE CLEMENT, LE MISERICORDIEUX

"Et on ne vous a donné que peu de connaissance"

[Saint Coran 17:85]

INTRODUCTION

Pendant longtemps, on a pensé que les phénomènes physiques sont invariants sous les transformations de conjugaison de charge, C, de parité, P et de renversement du sens du temps, T.

- L'opération de conjugaison de charge, C, transforme toute particule en son antiparticule. De ce fait, le signe de la charge électrique, de même que celui de tout autre nombre quantique additif de la particule (nombres leptoniques, baryonique, couleur, etc.), est inversé.
- La parité, P, n'est autre que l'opération d'inversion de l'espace, c'est-à-dire la symétrie par rapport à un point de l'espace. Elle inverse le signe des vecteurs de position ainsi que celui de toute grandeur vectorielle mais laisse invariantes les grandeurs pseudovectorielles (résultat du produit vectoriel de deux vrais vecteurs) telles que les moments cinétiques et, par cohérence, les spin.
- L'opération T consiste à inverser l'horloge, c'est-à-dire à remonter le temps. Inversant la coordonnée temporelle, T, va ainsi inverser les vecteurs impulsion et, de là, inverser aussi les moments cinétiques dont les spin.

Les physiciens pensaient, donc, que ces trois opérations constituaient de bonnes symétries de la nature, jusqu'en 1956 où la brisure de la symétrie de parité dans les interactions faibles fut prédite par T.D. Lee et C.N. Yang lorsqu'ils ont analysé les données expérimentales des désintégrations des mésons K chargés [1]. Cette brisure fut mise en évidence de façon irréfutable en 1957 par l'expérience de C.S. Wu qui examinait la désintégration bêta du noyau de cobalt. Les électrons résultants étaient préférentiellement émis dans la direction opposée à celle du moment magnétique du noyau de cobalt [2]. Une année plus tard (expérience de M. Goldhaber et son groupe [3]), on observait que les neutrinos émis lors de désintégrations faibles étaient toujours gauches (d'hélicité négative) alors que les antineutrinos étaient toujours droits (d'hélicité positive). Ceci est une autre évidence de la brisure de la symétrie P par les interactions faibles: dans une théorie invariante sous P, les deux états d'hélicité participeraient d'égale façon à l'interaction.

La théorie de l'interaction faible de E. Fermi (1934), modifiée par R.Feynman et M. Gell-Mann (1958) pour tenir compte de la brisure de la parité (théorie $V - A$) et améliorée par S. Glashow (1961) — qui proposa l'existence de trois bosons médiateurs W^+, W^- et Z — est en bon accord avec toutes ces observations expérimentales.

L'interaction faible brise aussi la symétrie C. On peut le voir dans la désintègration du π^+ en un positron et un neutrino. Ce dernier est toujours gauche. L'opération C transforme cet évènement en un π^- se désintègrant en un électron et un antineutrino gauche, ce qui n'est pas observé (l'antineutrino étant toujours droit).

Notons que C, P et T sont toujours de bonnes symétries en physique classique ainsi que dans les interactions fortes et électromagnétiques.

Peu après ces découvertes, Landau introduisit la combinaison CP comme nouvelle symétrie conservée. Il avait aussi observé que l'invariance CP avait comme conséquence la réalité des coefficients des lagrangiens des interactions faibles. Cependant, en 1964, le groupe de J.W. Cronin et V.L. Fitch observa la brisure de la symétrie CP dans les désintégrations des mésons K [4]. On savait avant cette expérience que les mésons de courte durée de vie, K_S^0, avaient surtout des modes de désintégration en deux pions (qui représentent un état propre de CP avec la valeur propre $+1$). Les mésons de longue durée de vie,

K_L^0, se désintègrent surtout en trois pions (représentant un état propre de CP avec la valeur propre -1) et en $\pi l \nu_l$. L'invariance CP expliquait parfaitement l'existence d'un seul mode par espèce — à savoir $\pi\pi$ pour K_S^0 et $\pi\pi\pi$ pour K_L^0 — car K_L^0 et K_S^0 étaient considérés comme états propres de CP avec les valeurs propres respectives -1 et $+1$. Mais dans l'expérience citée plus haut on observa que le méson K_L^0 se désintégrait en deux pions avec un taux d'embranchement de 3 sur 1000. Donc l'état propre de masse (ou état physique) K_L^0 n'est pas un état propre de CP. Ceci prouve que les interactions faibles brisent la symétrie CP.

D'autre part, le théorème CPT — selon lequel toute théorie quantique des champs locale et invariante de Lorentz est invariante sous la transformation CPT — implique que lorsqu'il y a brisure de CP il y a par conséquent brisure de T. Cette dernière a été mise en évidence, dans les interactions faibles, dans l'expérience CPLEAR au CERN [5], où la probabilité d'un $\overline{K^0}$ d'osciller en K^0 est légèrement supérieure (d'environ 6 sur 1000) à celle de l'oscillation inverse.

Depuis la découverte de la brisure CP dans le système des kaons, les physiciens sont toujours à la recherche d'autres manifestations de ce phénomène aussi bien en théorie qu'en expérience. De plus en plus de machines sont construites telles que les détecteurs BaBar (du PEP II au SLAC à Stanford, USA) et Belle (du KEK au Japon) qui viennent de fournir (2002) la première évidence expérimentale de la brisure CP dans les mésons B [6], car ce phénomène restait jusque là confiné au système des kaons neutres. On y reviendra un peu plus loin.

Pendant que la brisure CP restait mystérieuse, la physique des particules découvrait son modèle standard en 1968. Ce dernier, à son départ, ne considérait que deux générations de quarks. Kobayashi et Maskawa (1973) réalisaient qu'avec seulement deux générations, il était impossible de trouver dans le lagrangien une phase complexe nécessaire pour la brisure CP. Ils introduisirent, donc, la troisième génération, et montrèrent que la théorie avec trois générations avait une phase brisant CP. Comme il y a eu la découverte subséquente du quark b (en 1977 à Fermilab, Batavia, USA), ce modèle est devenu *le modèle standard de la brisure CP*. Presque deux décennies plus tard, en 1995, la découverte du quark t au Tevatron de Fermilab, donnera plus de crédibilité à ce modèle.

Il y a eu beaucoup de généralisations et d'extensions du modèle standard pour incorporer la brisure CP, mais on ne peut pas prétendre avoir trouvé une vraie explication à ce phénomène. On a développé une image assez claire sur les différentes façons de générer la brisure CP dans une théorie donnée, mais on est incapable de décider lequel de ces mécanismes est la source, ou encore, la source dominante de ce phénomène. La brisure CP demeure l'un des problèmes les plus importants de la physique des particules.

1 LE MODÈLE STANDARD

Le modèle standard de l'interaction électrofaible est un modèle de la théorie des champs quantiques qui décrit et unifie les interactions faibles et électromagnétiques (un survol des éléments du modèle standard est donné dans Réfs. [7, 8, 9, 10, 11]). C'est une théorie de jauge locale (où l'interaction est dictée par une invariance de jauge locale, i.e. que le lagrangien est invariant sous une transformation de phase dépendant des coordonnées), basée sur le groupe de symétrie $SU(2)_L \times U(1)_Y$. Les champs des particules (leptons et quarks) sont représentés en doublets gauches (i.e. d'hélicité négative ou L, de l'anglais *Left-Handed*) et en singulets droits (d'hélicité positive ou R de l'anglais *Right-Handed*):

Doublets de leptons gauches:

$$\begin{pmatrix} \nu_e \\ e^- \end{pmatrix}_L \quad \begin{pmatrix} \nu_\mu \\ \mu^- \end{pmatrix}_L \quad \begin{pmatrix} \nu_\tau \\ \tau^- \end{pmatrix}_L \tag{1}$$

Singulets de leptons droits:

$$e_R^- , \quad \mu_R^- , \quad \tau_R^- \tag{2}$$

Doublets de quarks gauches:

$$\begin{pmatrix} u \\ d \end{pmatrix}_L \quad \begin{pmatrix} c \\ s \end{pmatrix}_L \quad \begin{pmatrix} t \\ b \end{pmatrix}_L \tag{3}$$

Singulets de quarks droits:

$$u_R , \quad c_R , \quad t_R , \quad d_R , \quad s_R , \quad b_R \tag{4}$$

Les opérateurs de projection sur les états L et R sont donnés par

$$\gamma_L = \frac{1-\gamma_5}{2} \quad , \quad \gamma_R = \frac{1+\gamma_5}{2} \tag{5}$$

et pour tout champ X on a

$$X_L = \gamma_L X \quad , \quad X_R = \gamma_R X \ . \tag{6}$$

La théorie des groupes nous apprend qu'à chaque groupe de symétrie correspond un nombre bien déterminé de générateurs auquel correspond un nombre égal de bosons de jauge. Par exemple, l'électrodynamique quantique (ou QED de l'anglais *quantum electrodynamics*) est une théorie de jauge locale basée sur le groupe de symétrie $U(1)_{em}$. Ce groupe a un seul générateur qui est l'opérateur de charge électrique Q, auquel correspond un seul champ de jauge A_μ qui est le champ électromagnétique (le photon).

Dans le modèle standard, le groupe $U(1)$ introduit un seul boson de jauge, B_μ, et le groupe $SU(2)$ en introduit trois, W_μ^j, $j = 1, 2, 3$. Ces bosons sont sans masse alors que les bosons médiateurs de l'interaction faible doivent avoir des masses assez élevées pour expliquer la faiblesse de cette interaction. Donc pour générer des masses à ces bosons S. Weinberg et A. Salam ont fait appel au mécanisme de Higgs ou mécanisme de brisure spontanée de la symétrie (BSS) dans lequel on donne une masse à un boson de jauge en introduisant des champs annexes possédant une symétrie brisée (i.e. des états d'équilibre non symétriques). Grâce à ce mécanisme, partant de quatre bosons non massifs (cités plus haut), on se retrouve avec trois bosons massifs W^+, W^- et Z et un photon, γ, sans masse.

Ces médiateurs de l'intéraction électrofaible sont des combinaisons linéaires des quatre bosons de jauge cités plus haut:

$$\begin{aligned} W^{\pm}_{\mu} &= \frac{1}{\sqrt{2}}(W^1_{\mu} \pm iW^2_{\mu}) \\ Z_{\mu} &= -\sin\theta_W B_{\mu} + \cos\theta_W W^3_{\mu} \\ \gamma_{\mu} &= \cos\theta_W B_{\mu} + \sin\theta_W W^3_{\mu} \end{aligned} \tag{7}$$

et les masses que gagnent $W^{\pm}$ et Z par le mécanisme de Higgs sont

$$M_W = \frac{37.3 \text{ GeV}}{\sin\theta_W} \quad , \quad M_Z = \frac{M_W}{\cos\theta_W} \ . \tag{8}$$

Ici, θ_W, appelé angle de Weinberg, est un paramètre libre de la théorie. De plus, on se retrouve avec un nouveau boson, de spin et de charge électrique nuls, *le boson de Higgs.*

Le modèle standard a eu un succès remarquable dans la description des phénomènes à faible énergie. Il a prédit avec succès les propriétés des bosons $W^{\pm}$ et Z. En effet, les courants neutres (interactions par échange du Z) furent découverts expérimentalement au CERN en 1973. Puis en 1983, toujours au CERN, il y a eu la découverte des bosons $W^{\pm}$ et Z et on mesura $\sin^2\theta_W = 0.232$.

Il reste que ce modèle comporte des lacunes. L'arbitrarité des masses des fermions (générées elles aussi par BSS) et de leurs couplages, et l'arbitrarité aussi de la masse du boson de Higgs, non observé jusqu'à ce jour, en sont les principales. La difficulté d'observation du boson de Higgs est due à son faible couplage aux fermions (de l'ordre de $(m_f/M_W)e$). Les limites actuelles pour sa masse sont 100 GeV $< M_H <$ 1 TeV.

2 MATRICE DE KOBAYASHI-MASKAWA

Plus tôt dans l'introduction, il a été mentionné que la matrice de Kobayashi-Maskawa est devenue le modèle standard de la brisure de CP. Voyons, de brève façon, comment se construit cette matrice. (Plus de détails se trouvent dans Réfs. [8, 12, 13, 14].)

Soit le lagrangien d'interaction quarks-Higgs

$$\mathcal{L}(q,H) = \sum_{j,k=1}^{3} \left\{ Y_{jk}\overline{Q_j}_L H^C q_{k_R} + Y'_{jk}\overline{Q_j}_L H q'_{k_R} + h.c \right\} , \tag{9}$$

où $Y^{(')}_{jk}$ sont des constantes de couplage et "*h.c.*" signifie l'hermétien conjugué de l'expression qui précède. Ici, Q_{j_L} est un doublet de quarks gauches [éq. (3)], q_{k_R} et q'_{k_R} sont des singulets de quarks droits, respectivement "*up*" et "*down*" [éq. (4)] et H est le doublet de Higgs défini par

$$H \equiv \begin{pmatrix} \phi^{(+)} \\ \phi^{(0)} \end{pmatrix} \tag{10}$$

avec

$$\begin{aligned} \phi^{(+)} &= \frac{1}{\sqrt{2}}(\phi_1 + i\phi_2) \\ \phi^{(0)} &= \frac{1}{\sqrt{2}}(\phi_0 + i\phi_3) \ . \end{aligned} \tag{11}$$

Trois de ces scalaires , ϕ_1, ϕ_2 et ϕ_3, seront "mangés" (mécanisme de Higgs) par les bosons de jauge qui deviendront massifs et ne restera qu'un seul Higgs physique, ϕ_0. H^C est le C-conjugué de H.

Sous la brisure spontanée de la symétrie on a

$$\phi_0 \longrightarrow \phi_0 + v \ , \tag{12}$$

où v est la valeur dans le vide de ϕ_0. Ainsi, le lagrangien de l'éq. (9) devient

$$\mathcal{L}(q,H) \longrightarrow -\sum_{j,k=1}^{3} \left\{ m_{jk}\overline{q_j}_L q_{k_R} + m'_{jk}\overline{q'_j}_L q'_{k_R} + h.c. \right\} \left(1 + \frac{1}{v}\phi_0\right) \tag{13}$$

où

$$m_{jk} = -\frac{v}{\sqrt{2}} Y_{jk} \quad , \quad m'_{jk} = -\frac{v}{\sqrt{2}} Y'_{jk} \tag{14}$$

sont des éléments des *matrices de masse* des quarks, m pour les quarks "*up*" et m' pour les quarks "*down*". Ces matrices ne sont pas diagonales, donc les champs $q_j^{(')}$ ne sont pas des champs physiques. De plus, dans cette base non physique, m et m' sont complexes. Elles renferment des phases brisant CP.

Maintenant, on va diagonaliser les matrices m et m' à l'aide de matrices unitaires U_L, U_R, U'_L et U'_R, comme suit:

$$\begin{aligned} U_L m U_R^\dagger &= D \equiv \begin{pmatrix} m_u & 0 & 0 \\ 0 & m_c & 0 \\ 0 & 0 & m_t \end{pmatrix} \ , \\ U'_L m' U_R'^\dagger &= D' \equiv \begin{pmatrix} m_d & 0 & 0 \\ 0 & m_s & 0 \\ 0 & 0 & m_b \end{pmatrix} \ . \end{aligned} \tag{15}$$

Les m_i, $i = u, c, t, d, s, b$, sont les masses des quarks, correspondant à la base physique (contenant les vecteurs propres de m et m'):

$$q^{ph}_{L(R)} \equiv \begin{pmatrix} u \\ c \\ t \end{pmatrix}_{L(R)} = U_{L(R)}\, \tilde{q}_{L(R)} \equiv U_{L(R)} \begin{pmatrix} \tilde{u} \\ \tilde{c} \\ \tilde{t} \end{pmatrix}_{L(R)}$$

$$q'^{ph}_{L(R)} \equiv \begin{pmatrix} d \\ s \\ b \end{pmatrix}_{L(R)} = U'_{L(R)}\,\tilde{q}'_{L(R)} \equiv U'_{L(R)} \begin{pmatrix} \tilde{d} \\ \tilde{s} \\ \tilde{b} \end{pmatrix}_{L(R)} . \qquad (16)$$

$\tilde{u}, \tilde{c}$,... etc. sont les champs de quarks non-physiques, ceux participant au lagrangien de l'éq. (9) qui va s'écrire dans la base physique

$$\mathcal{L}^{ph}(q, H) = -\left(1 + \frac{1}{v}\phi_0\right)\left[m_u\bar{u}u + m_c\bar{c}c + m_t\bar{t}t + m_d\bar{d}d + m_s\bar{s}s + m_b\bar{b}b\right] . \qquad (17)$$

Ce dernier terme conserve parfaitement la symétrie CP. On va voir que la brisure de CP a éte transférée aux courants chargés.

Le terme de courant chargé de la partie hadronique de la densité du lagrangien du modèle standard est

$$\mathcal{L}_C(q) = \sum_{j=1}^{3}\left[-\frac{g_2}{\sqrt{2}}\right] W^-_\mu \overline{\tilde{q}}_{j\,L}\gamma^\mu \tilde{q}'_{j_L} + h.c. \qquad (18)$$

et ce, dans la base non-physique des quarks. g_2 est la constante de couplage faible aux courants chargés. Elle est reliée à la charge électrique par la relation

$$e = g_2 \sin\theta_W \ , \qquad (19)$$

θ_W étant l'angle de Weinberg. L'expression de l'éq. (18) dans la base physique donne, en utilisant l'éq. (16)

$$\begin{aligned} \mathcal{L}^{ph}_C(q) &= \left[-\frac{g_2}{\sqrt{2}}\right] W^-_\mu \overline{q^{ph}_L}\gamma^\mu U_L U'^\dagger_L q'^{ph}_L + h.c. \\ &\equiv \left[-\frac{g_2}{\sqrt{2}}\right] W^-_\mu \overline{q^{ph}_L}\gamma^\mu V_{KM} q'^{ph}_L + h.c., \end{aligned} \qquad (20)$$

où

$$V_{KM} = \begin{pmatrix} V_{ud} & V_{us} & V_{ub} \\ V_{cd} & V_{cs} & V_{cb} \\ V_{td} & V_{ts} & V_{tb} \end{pmatrix} \qquad (21)$$

est la matrice de mélange des quarks connue sous le nom de matrice de Kobayashi-Maskawa (ses inventeurs) [15].

Il y a donc mélange de saveurs dans l'interaction à courant chargé et on voit que ce courant chargé est toujours gauche dans le modèle standard. Voyons ce qui se passe dans les couplages aux courants neutres.

Le terme de courant neutre de la partie hadronique de la densité du lagrangien du modèle standard est, dans la base non-physique des quarks,

$$\begin{aligned}\mathcal{L}_N(q) = \left[-\frac{g_2}{\sqrt{2}\cos\theta_W}\right] \; Z^0_\mu \; & \left[\overline{\tilde{q}_L}\,\gamma^\mu\left(\frac{1}{2}-\frac{2}{3}\sin^2\theta_W\right)\tilde{q}_L\right. \\ & +\overline{\tilde{q}_R}\,\gamma^\mu\left(-\frac{2}{3}\sin^2\theta_W\right)\tilde{q}_R \\ & +\overline{\tilde{q}'_L}\,\gamma^\mu\left(-\frac{1}{2}+\frac{1}{3}\sin^2\theta_W\right)\tilde{q}'_L \\ & \left.+\overline{\tilde{q}'_R}\,\gamma^\mu\left(\frac{1}{3}\sin^2\theta_W\right)\tilde{q}'_R\right] \;. \end{aligned} \tag{22}$$

En utilisant les relations de l'éq. (16), ce terme donne, sachant que les quatre matrices $U^{(')}_{L(R)}$ sont unitaires,

$$\begin{aligned}\mathcal{L}^{ph}_N(q) = \left[-\frac{g_2}{\sqrt{2}\cos\theta_W}\right] \; Z^0_\mu \; & \left[\overline{q^{ph}_L}\,\gamma^\mu\left(\frac{1}{2}-\frac{2}{3}\sin^2\theta_W\right)q^{ph}_L\right. \\ & +\overline{q^{ph}_R}\,\gamma^\mu\left(-\frac{2}{3}\sin^2\theta_W\right)q^{ph}_R \\ & +\overline{q'^{ph}_L}\,\gamma^\mu\left(-\frac{1}{2}+\frac{1}{3}\sin^2\theta_W\right)q'^{ph}_L \\ & \left.+\overline{q'^{ph}_R}\,\gamma^\mu\left(\frac{1}{3}\sin^2\theta_W\right)q'^{ph}_R\right] \;. \end{aligned} \tag{23}$$

De là on constate que, dans le modèle standard, il n'existe pas de courants neutres avec changement de saveur (FCNC ou "*flavour-changing neutral currents*"). Autrement dit, le boson Z ne couple que les quarks de même saveur. La matrice de Kobayashi-Maskawa n'intervient que dans les courants chargés. De plus, les courants neutres peuvent être droits ou gauches, contrairement aux courants chargés qui sont toujours gauches.

Retournons à la matrice de Kobayashi-Maskawa dont les élements doivent être déterminés expérimentalement via les couplages $W^\pm\overline{q}q'$. La théorie des groupes nous apprend qu'une matrice complexe et unitaire $N \times N$ est décrite par $\frac{1}{2}N(N-1)$ paramètres réels indépendants et $\frac{1}{2}N(N+1)$ phases complexes. Étant donné que V_{KM} est une matrice complexe et unitaire 3×3, elle doit donc comprendre trois paramètres réels indépendants et six phases complexes. Cependant, ces six phases n'ont pas toutes un sens physique. Sachant que l'interaction de l'éq. (20) est invariante sous transformation de phases, on peut "absorber" cinq des six phases complexes de V_{KM}, en réajustant les phases des composantes droites des quarks. Cependant, une phase reste redondante. *C'est cette phase unique de V_{KM} qui cause la brisure CP dans le modèle standard.* Ajoutons enfin que le changement de base (passage des quarks non-physiques aux quarks physiques, décrit par l'éq. (16)) a transféré la brisure CP des couplages quarks-Higgs aux couplages quarks-boson W.

Il existe plusieurs paramétrisations de la matrice V_{KM} mais citons ici la plus pratique, à savoir celle de Wolfenstein [16], qui met en relief l'ordre de grandeur de chaque élément et la phase unique brisant CP représentée par le paramètre η (A, ρ et λ sont les trois paramètres réels indépendants mentionnés plus haut). Dans cette paramétrisation, la matrice V_{KM} s'écrit approximativement

$$V_{KM}^{W} \simeq \begin{pmatrix} 1-\frac{1}{2}\lambda^2 & \lambda & A\lambda^3(\rho - i\eta) \\ -\lambda(1+iA^2\lambda^4\eta) & 1-\frac{1}{2}\lambda^2 & A\lambda^2 \\ A\lambda^3(1-\rho-i\eta) & -A\lambda^2 & 1 \end{pmatrix} , \qquad (24)$$

où les éléments présentent une hiérarchie en termes du paramètre λ (sinus de l'angle de Cabibbo). Notons qu'à l'ordre $0(\lambda^3)$, il y a seulement deux élements de la matrice V_{KM} qui ont une partie imaginaire non nulle, V_{ub} et V_{td}. Les mesures expérimentales ont permis de trouver [17]

$$\begin{aligned} A &= 0.819 \pm 0.035 \ , \\ \lambda &= 0.2196 \pm 0.0023 \ , \\ \sqrt{\rho^2+\eta^2} &= 0.423 \pm 0.064 \ . \end{aligned} \qquad (25)$$

Il y a trois autres paramètres de la matrice V_{KM} qui sont souvent cités dans

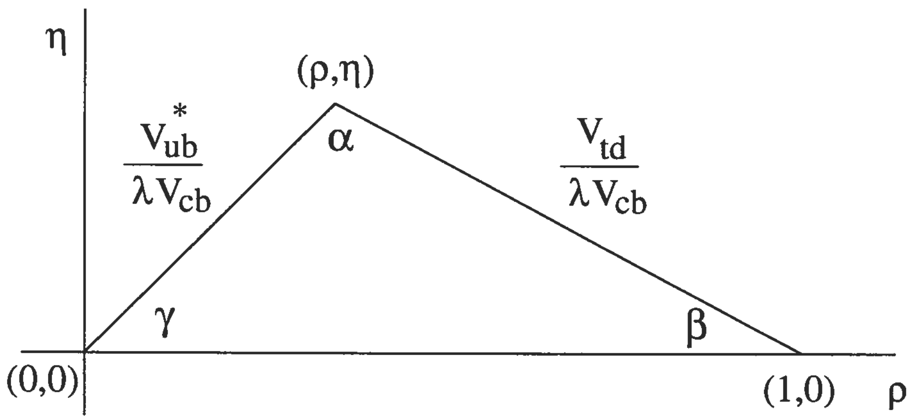

Figure 1: *Le triangle d'unitarité.*

le problème de la brisure CP. Ce sont les trois angles du triangle d'unitarité. Ce dernier n'est autre qu'une représentation géométrique dans le plan complexe de la condition d'unitarité de la matrice V_{KM}. Soit

$$V_{ud}V_{ub}^* + V_{cd}V_{cb}^* + V_{td}V_{tb}^* = 0 \ , \tag{26}$$

qui peut être réecrite en utilisant l'éq. (24) comme

$$\frac{V_{ub}^*}{\lambda V_{cb}} + \frac{V_{td}}{\lambda V_{cb}} = 1 \ . \tag{27}$$

Cette dernière équation représente un triangle dans le plan complexe (i.e. le plan ρ–η), illustré dans la Fig. 1. C'est le triangle d'unitarité. Ainsi , les valeurs permises de ρ et η se traduisent en les formes permises du triangle d'unitarité.

Les angles α, β et γ de ce triangle sont directement reliés à la phase complexe de la matrice V_{KM}. En effet,

$$\alpha \ \equiv \ arg\left(-\frac{V_{td}V_{tb}^*}{V_{ud}V_{ub}^*}\right) \simeq \arctan\left[\frac{-\eta}{1-\rho}\right] - \arctan\left[\frac{\eta}{\rho}\right] \ ,$$

$$\beta \equiv arg\left(-\frac{V_{cd}V_{cb}^*}{V_{td}V_{tb}^*}\right) \simeq \pi - \arctan\left[\frac{-\eta}{1-\rho}\right] ,$$
$$\gamma \equiv arg\left(-\frac{V_{ud}V_{ub}^*}{V_{cd}V_{cb}^*}\right) \simeq \arctan\left[\frac{\eta}{\rho}\right] - \pi . \tag{28}$$

Alors, il y a brisure CP (i.e., $\eta \neq 0$) lorsque l'aire du triangle d'unitarité n'est pas nulle, autrement dit, lorsque les trois angles α, β et γ sont non nuls. (Dans le modèle standard, leur somme géométrique vaut π. Ceci est un test du MS.)

De récentes mesures expérimentales (2002) dans les nouveaux détecteurs Babar et Belle ont permis de mesurer $\sin 2\beta = 0.78 \pm 0.08$ dans les désintégrations des mésons B [6]. Ceci donne la première évidence expérimentale de la brisure CP hors du système des kaons.

En fait, la mesure des angles α, β et γ est la raison d'être des "usines" à mésons B pour trouver la moindre contradiction avec le modèle standard. Le modèle standard explique bien le phénomène de la brisure CP observé dans les désintégrations des kaons. Il présente un mécanisme "posé à la main", à l'aide de la matrice de Kobayashi-Maskawa, mais il ne nous apprend rien sur l'origine de ce phénomène. C'est pour cela qu'il faut chercher encore plus "d'indices" dans tous les endroits possibles en particulier dans les désintégrations des hadrons à saveur de beauté, à savoir des mésons B et des hypérons Λ_b.

3 BRISURE CP DANS LES MÉSONS B

Le modèle standard prédit de grands effets de brisure CP dans les désintégrations des hadrons faisant intervenir les quarks b (des détails se trouvent dans Réfs. [7, 9, 12, 18, 19, 20, 21, 22]). Ces derniers, appartenant à la troisième famille, ont beaucoup de modes de désintégration en membres de familles inférieures. Ainsi, les mésons B, comme les hypérons Λ_b, se désintègrent en une multitude d'états finaux où des effets d'asymétries CP peuvent être mesurés.

Une étude quantitative de ces asymétries conduirait soit à la conclusion que la phase de la matrice V_{KM} est la seule source de brisure CP soit à des bases expérimentales pour de nouveaux mécanismes de brisure CP, amenant l'étude au-delà du modèle standard. Donc l'étude des asymétries CP dans les hadrons beaux (dont l'un des quarks constituants est le quark b) devrait répondre à la question: est-ce que l'unique phase complexe de la matrice V_{KM} est la seule source de brisure CP ? En outre, cette étude permettrait aussi une mesure plus exacte des paramètres de la matrice V_{KM}.

La présente thèse est une contribution à cette étude. Elle concerne la recherche d'asymétries de brisure CP via les produits triples. On va y revenir dans la section suivante.

Pour le moment, intéressons-nous à un autre type de signaux de brisures CP, les plus étudiés en littérature, soient les asymétries des taux de désintégration. Ces asymétries peuvent avoir lieu dans les désintégrations des mésons B neutres ou chargés. Dans le cas des B neutres des effets de mélange (on va voir plus loin ce qu'on entend par effet de mélange) peuvent se manifester et contribuer au signal. Si aucune contribution des effets de mélange n'a lieu (ce qui est toujours le cas pour les B chargés) on parle d'asymétries de brisure CP directe. Il peut aussi y avoir brisure CP due uniquement aux effets de mélange ou encore combinaison de brisure directe et d'effets de mélange.

3.1 Brisure CP indirecte dans les mésons B

Dans le système des mésons B neutres, les états physiques (états propres de masse et de l'interaction faible) sont des superpositions des états conjugués de CP, B^0 et $\bar{B}^0$. Ils s'écrivent

$$\begin{aligned} |B_L\rangle &= p|B^0\rangle + q|\bar{B}^0\rangle \\ |B_H\rangle &= p|B^0\rangle - q|\bar{B}^0\rangle \ . \end{aligned} \tag{29}$$

Les indices L et H sont pour, respectivement, "*light*" (léger) et "*heavy*" (lourd). En fait, pour les mésons B, c'est la différence de masse $\Delta m_B = m_{B_H} - m_{B_L}$ qui domine la physique, alors que les durées de vie des deux états physiques sont presque les mêmes. Ce "mélange" de B^0 et $\bar{B}^0$ est dû au fait que le nombre quantique de beauté, b, n'est pas conservé par les interactions faibles. De plus, la brisure de CP implique que $p \neq q$, car si CP était une bonne symétrie, on aurait $p = q$ et $|B_L\rangle$ et $|B_H\rangle$ seraient des états propres de CP avec les valeurs propres respectives $+1$ et -1. La différence entre p et q caractérise la brisure CP due au mélange B^0–$\bar{B}^0$. D'autre part, pour le système des mésons B, p et q ne diffèrent que d'une phase, ϕ_M, phase faible[1] de ce mélange. On a

$$q/p \approx e^{-2i\phi_M} \ . \tag{30}$$

Dans le modèle standard, les prédictions les plus importantes de signaux de brisure CP qu'on trouve dans la littérature ont lieu dans les désintégrations des mésons B neutres en états propres de CP. Ces signaux sont directement liés aux angles du triangle d'unitarité. Voyons brièvement comment.

Pour une désintégration $B^0 \to f_{CP}$, l'asymétrie CP du taux de désintégration est définie par

$$a_{f_{CP}}(t) = \frac{\Gamma(B^0(t) \to f_{CP}) - \Gamma(\overline{B}^0(t) \to f_{CP})}{\Gamma(B^0(t) \to f_{CP}) + \Gamma(\overline{B}^0(t) \to f_{CP})} \ , \tag{31}$$

où $B^0(t)$ (resp. $\overline{B}^0(t)$) est un méson B neutre à l'instant t qui était un B^0 (resp. un $\bar{B}^0$) à l'instant $t = 0$. On démontre que

$$a_{f_{CP}}(t) = \frac{(1 - |r_{f_{CP}}|^2)\cos(\Delta m_B t) - 2Im r_{f_{CP}} \sin(\Delta m_B t)}{1 + |r_{f_{CP}}|^2} \ , \tag{32}$$

[1]On veut dire par phases faibles les phases présentes dans les amplitudes de désintégration qui ont pour origine la phase de la matrice de Kobayashi-Maskawa. Elles sont toutes reliées a cette dernière et changent de signe entre B^0 et $\bar{B}^0$. Les phases fortes, par contre, proviennent des interactions fortes entre états finaux et sont les mêmes pour B^0 et $\bar{B}^0$.

où Δm_B est la difference de masse mentionnée plus haut et le paramètre $r_{f_{CP}}$ est défini par

$$r_{f_{CP}} \equiv \frac{q}{p}\frac{\bar{A}}{A} , \tag{33}$$

avec A représentant l'amplitude de désintégration à $t = 0$ de B^0 en f_{CP} et $\bar{A}$ l'amplitude du processus CP-conjugué.

Dans le cas de brisure CP due seulement au mélange, on a, $|r_{f_{CP}}| = 1$ car A et $\bar{A}$ ont, dans ce cas, une même phase faible, ϕ_D,

$$\bar{A}/A = e^{-2i\phi_D} \tag{34}$$

et l'éq. (32) se simplifie en

$$a_{f_{CP}}(t) = -Imr_{f_{CP}} \sin(\Delta m_B t) \tag{35}$$

avec, en utilisant les éqs. (30) et (34)

$$Imr_{f_{CP}} = -\sin 2(\phi_M + \phi_D) . \tag{36}$$

On voit que l'asymétrie CP due au mélange [éq. (35)] *oscille dans le temps avec, comme fréquence, la différence entre les masses des B^0 physiques (Δm_B) et une amplitude ($Imr_{f_{CP}}$) qui est une pure fonction des paramètres de V_{KM}.*

Voyons maintenant des exemples typiques de détermination des angles, α, β et γ, du triangle d'unitarité à partir d'asymétries du type donné dans l'éq. (35).

On considère d'abord la désintégration $B_d^0 \to \psi K_S$. Elle correspond au sous-processus $b \to c\bar{c}s$. Comme ce sont les couplages b-c et s-$\bar{c}$ qui sont mis en jeu, la phase faible de désintégration [éq. (34)] n'est autre que

$$-2\phi_D = arg\left(\frac{\bar{A}}{A}\right) = arg\left(\frac{V_{cb}V_{cs}^*}{V_{cb}^*V_{cs}}\right) . \tag{37}$$

D'autre part, on démontre que la phase faible du mélange $B_d^0 - \overline{B_d^0}$ [éq. (30)] est

$$arg\left(\frac{q}{p}\right)_{B_d^0} = arg\left(\frac{V_{td}V_{tb}^*}{V_{tb}V_{td}^*}\right) . \tag{38}$$

Mais il y a aussi le mélange $K^0 - \bar{K}^0$ dans K_S qui correspond à la phase faible

$$arg\left(\frac{q}{p}\right)_{K_S} = arg\left(\frac{V_{cs}V_{cd}^*}{V_{cd}V_{cs}^*}\right) . \tag{39}$$

Donc, la phase totale de mélange est

$$-2\phi_M = arg\left\{\left(\frac{V_{td}V_{tb}^*}{V_{tb}V_{td}^*}\right)\left(\frac{V_{cs}V_{cd}^*}{V_{cd}V_{cs}^*}\right)\right\} . \tag{40}$$

Des éqs. (37) et (40), on tire le paramètre $r_{f_{CP}}$ [éq. (36)] pour $f_{CP} = \psi K_S$:

$$arg\,(r_{\psi K_S}) = -2(\phi_D + \phi_M) = -2arg\left(\frac{V_{cd}V_{cb}^*}{V_{td}V_{tb}^*}\right) = 2\beta . \tag{41}$$

La dernière égalité provenant de l'éq. (28). L'amplitude de l'asymétrie CP dans le mode $B_d^0 \to \psi K_S$ est directement liée à l'angle β du triangle d'unitarité, soit

$$Im\,(r_{\psi K_S}) = \sin(2\beta) . \tag{42}$$

De façon analogue, on peut mesurer $\sin(2\alpha)$ dans la désintégration $B_d^0 \to \pi^+\pi^-$ et $\sin(2\gamma)$ dans $B_s^0 \to \rho K_S$.

On voit de ces trois exemples comment le Modèle Standard offre la possibilité de mesurer indépendamment les trois angles du triangle d'unitarité. L'angle β a été récemment mesuré dans les détecteurs BaBar et Belle [6] comme mentionné dans le chapitre précédent et la désintégration $B_d^0 \to \psi K_S$ fait partie des modes étudiés dans ces expériences. On a trouvé $\sin 2\beta = 0.78 \pm 0.08$.

Avant de passer au sujet principal de la thèse (qui est la brisure CP via les produits triples) jetons un coup d'œil sur la brisure CP directe dans les mésons B.

3.2 Brisure CP directe dans les mésons B

Pour une désintégration $B \to f$, sans effets de mélange, l'asymétrie CP directe est définie par

$$a_f = \frac{\Gamma(B \to f) - \Gamma(\bar{B} \to \bar{f})}{\Gamma(B \to f) + \Gamma(\bar{B} \to \bar{f})} , \tag{43}$$

où B peut être chargé ou neutre. Pour qu'une telle asymétrie ait lieu, il faut la contribution d'au moins deux amplitudes au taux de désintégration du mode, avec différentes phases faibles et différentes phases fortes, à savoir

$$\begin{aligned} A(B \to f) &= A_1 e^{i(\delta_1 + \phi_1)} + A_2 e^{i(\delta_2 + \phi_2)} \\ A(\overline{B} \to \overline{f}) &= A_1 e^{i(\delta_1 - \phi_1)} + A_2 e^{i(\delta_2 - \phi_2)} . \end{aligned} \tag{44}$$

où les ϕ_i sont les phases faibles et les δ_i sont les phases fortes. L'asymétrie [éq. (43)] va donc être

$$a_f \propto A_1 A_2 \sin(\delta_1 - \delta_2) \sin(\phi_1 - \phi_2) . \tag{45}$$

Ces asymétries sont aussi utilisées pour tester le modèle standard en mesurant les angles α, β et γ du triangle d'unitarité.

Le problème avec ce genre d'asymétries est que la phase forte relative, $(\delta_1 - \delta_2)$ [éq. (45)] doit absolument être non nulle pour que le signal de brisure CP soit détecté. De plus cette phase forte constitue souvent un paramètre inconnu. On va voir, dans ce qui suit, qu'il en est autrement pour les asymétries qui nous intéressent: celles dues aux produits triples.

4 PRODUITS TRIPLES ET BRISURE CP

Les asymétries CP via les produits triples ont été très peu étudiées comparativement aux signaux de brisure CP qui ont reçu le plus d'attention et qui sont les asymétries des taux de désintégration dans les mésons B mentionnées dans la section précédente. La présence de produits triples dans les amplitudes de désintégration est une manifestation directe de brisure de la symétrie T et, en tenant compte du théorème CPT, elle est aussi un signal de brisure de la symétrie CP.

On entend par produit triple tout scalaire $\vec{v}_1 \cdot (\vec{v}_2 \times \vec{v}_3)$ où chaque v_i est soit une impulsion soit un vecteur spin de l'une des particules participant à la désintégration. On avait vu, au début de ce chapitre, que l'opération T inversait le signe de ces vecteurs. Ainsi, les produits triples ont toujours leurs signes inversés par l'opération T, de sorte que leur présence dans une amplitude de désintégration indique que cette dernière n'est pas un invariant de la symétrie T.

On définit pour les produits triples l'asymétrie

$$A_T = \frac{\Gamma\left(\vec{v}_1 \cdot (\vec{v}_2 \times \vec{v}_3) > 0\right) - \Gamma\left(\vec{v}_1 \cdot (\vec{v}_2 \times \vec{v}_3) < 0\right)}{\Gamma\left(\vec{v}_1 \cdot (\vec{v}_2 \times \vec{v}_3) > 0\right) + \Gamma\left(\vec{v}_1 \cdot (\vec{v}_2 \times \vec{v}_3) < 0\right)}\,, \tag{46}$$

où Γ est le taux de désintégration du processus étudié. Cette asymétrie indique que le taux de désintégration dans la région de l'espace où le produit triple est positif diffère du taux de désintégration dans la région de l'espace où le produit triple est négatif. Mais A_T ne représente pas en soit la brisure de T car T ne fait pas qu'inverser les signes des spins et des impulsions mais inverse aussi états finaux et états initiaux. Comme on n'inverse pas les états initiaux et finaux dans A_T, on peut donc avoir $A_T \neq 0$ sans avoir de brisure T. On parle dans ce cas d'un faux signal de brisure T (CP). On va voir plus loin comment cela se produit et comment discerner les vrais signaux des faux signaux. Comme A_T n'indique pas si le signal est vrai ou faux, on ne l'appelle pas encore asymétrie de brisure T mais plutôt *asymétrie T impaire*.

Supposons qu'on ait une désintégration $B \to f$ dominée par les deux amplitudes de l'éq. (44) dont l'interférence contient un produit triple. Alors l'asymétrie T impaire va être

$$A_T \propto \cos(\delta_1 - \delta_2)\sin(\phi_1 - \phi_2) + \sin(\delta_1 - \delta_2)\cos(\phi_1 - \phi_2) \ . \tag{47}$$

On voit bien dans l'éq. (47) que A_T peut être non nulle même si la phase faible $(\phi_1 - \phi_2)$ est nulle. C'est le faux signal dont on parlait plus haut. Il est dû à la présence des phases fortes [terme $\sin(\delta_1 - \delta_2)$]. Pour remédier à cela on définit *l'asymétrie de brisure T* comme suit:

$$\mathcal{A}_T = \frac{1}{2}(A_T \pm \overline{A}_T) \ , \tag{48}$$

où $\overline{A}_T$ est l'asymétrie T impaire du processus CP-conjugué $\overline{B} \to \overline{f}$. Le signe qu'on met entre A_T et $\overline{A}_T$ dépend du type de PT; on met un signe (+) lorsque le nombre d'impulsions y est impair et un signe (−) lorsqu'il est pair. Ceci provient du fait que CP inverse les impulsions, donc inverse le signe du premier type de PT. $\mathcal{A}_T$ donne

$$\mathcal{A}_T \propto \cos(\delta_1 - \delta_2)\sin(\phi_1 - \phi_2) \ . \tag{49}$$

Il est clair que $\mathcal{A}_T$ n'est non nulle que lorsque la phase faible $(\phi_1 - \phi_2)$ est non nulle. $\mathcal{A}_T$ est donc un vrai signal de brisure T (CP).

D'autre part, l'avantage de cette classe d'asymétries par rapport aux asymétries des taux de désintégration [éqs. (43) et (45)] est que la présence de phases fortes n'est pas nécessaire pour qu'elle ait lieu. Au contraire, le signal est maximal lorsque la phase forte est nulle [dans l'éq. (49), $\mathcal{A}_T$ est maximale pour $\delta_1 - \delta_2 = 0$].

Ajoutons qu'on peut définir l'asymétrie suivante (le signe qu'on met entre A_T et $\overline{A}_T$ est l'opposé de celui qu'on met dans $\mathcal{A}_T$):

$$\mathcal{A}_T^F = \frac{1}{2}(A_T \mp \overline{A}_T) \propto \sin(\delta_1 - \delta_2)\cos(\phi_1 - \phi_2) \ , \tag{50}$$

Dans les désintégrations où la phase faible est nulle, $\mathcal{A}_T^F$ n'est pas un vrai signal de brisure CP; c'est le faux signal dont on parlait au début de cette section. Il va permettre de mesurer la phase forte:

$$\mathcal{A}_T^F\left((\phi_1-\phi_2)=0\right)\propto\sin\left(\delta_1-\delta_2\right)\,. \tag{51}$$

Maintenant, voilà comment se manifeste un produit triple, $\vec{v}_1\cdot(\vec{v}_2\times\vec{v}_3)$, dans le carré d'une amplitude. Tout terme de la forme $\varepsilon_{\mu\nu\rho\sigma}p^\mu v_1^\nu v_2^\rho v_3^\sigma$, où $\varepsilon_{\mu\nu\rho\sigma}$ est le tenseur complètement antisymétrique à quatre dimensions et p^μ est la quadri-impulsion de l'une des particules participant à la désintégration, se réduit dans le référentiel propre de cette dernière au terme $m\vec{v}_1\cdot(\vec{v}_2\times\vec{v}_3)$, où m est la masse de la particule.

La recherche de signaux de brisure CP via les produits triples est un domaine très intéressant de la physique des particules. De plus en plus d'usines à mésons B se construisent à travers le monde , donnant la chance de mesurer de tels signaux, de tester la validité du modèle standard et en même temps de diriger la recherche vers la bonne physique au-delà du modèle standard. La présente thèse vient déposer une nouvelle brique dans la construction d'une idée claire sur l'origine du phénomène de la brisure CP. Les conséquences de la compréhension de ce phénomène s'étendront jusqu'en cosmologie, où on pourra expliquer la quasi-inexistence de l'antimatière dans notre monde.

Le premier article de cette thèse est consacré à l'étude d'asymétries de brisure T via les produits triples [du type donné dans l'éq. (48)] dans le processus inclusif $b\to\ su\bar{u}$ et ce, dans le modèle standard. Les applications au niveau hadroniques (mésons B et hypérons Λ_b) sont mentionnées. Trois amplitudes contribuent à la désintégration $b\to su\bar{u}$ donnant lieu à plusieurs produits triples dans l'amplitude totale au carré (plus précisément dans les termes d'interférence de celle-ci). Deux seulement de ces produits triples donnent des asymétries assez larges. L'un fait intervenir les spins des quarks u et $\bar{u}$ avec soit l'impulsion de u soit celle de $\bar{u}$ et l'autre produit triple fait intervenir le spin du quark b avec les impulsions des quarks u et s. Les asymétries trouvées sont d'environ 5% pour le premier produit triple et entre 1% et 3% pour le second. Les asymétries dues aux autres produits triples sont négligeables, entre autres ceux impliquant le spin du quark s (asymétrie réduite à cause de la suppression en $1/m_b$ dans l'amplitude). L'idée est que si l'asymétrie due à un produit triple est négligeable au niveau des quarks elle doit l'être aussi au niveau des hadrons, ce qui permet de tester le MS et de voir, s'il y a lieu, la manifestation d'une nouvelle physique.

La première asymétrie peut être mesurée dans des désintégrations $B \to V_1 V_2$, où V_1 et V_2 sont des mésons vecteurs. Par contre, comme la deuxième asymétrie implique le spin du quark b, elle ne peut pas être mesurée dans des désintégrations des mésons B; il faut utiliser des baryons Λ_b. Ceci est étudié dans les deuxième et troisième articles, qui constituent donc une suite logique au premier.

Dans le deuxième article, on reste dans le modèle standard et on calcule les asymétries de brisure T dues aux produits triples dans les désintégrations non charmées de Λ_b en utilisant le modèle de factorisation. On trouve une grande asymétrie (18%) pour le mode $\Lambda_b \to pK^-$ correspondant au mode inclusif $b \to su\bar{u}$. Ici la suppression en $1/m_b$, mentionnée plus haut, est compensée par le terme d'accroissement chiral ("*chiral enhancement*"), $m_K^2/(m_s + m_u)$. Pour le mode $\Lambda_b \to p\pi^-$, on trouve une asymétrie de l'ordre de 5% et elle est de l'ordre de 1% pour $\Lambda_b \to \Lambda\eta(\eta')$. Pour les modes $\Lambda_b \to FV$, où F est un baryon léger de spin 1/2 (un proton ou un Λ) et V un méson vecteur (un K^*, un ρ ou un ϕ), les asymétries sont très petites, valant au plus 1% pour les V polarisés transversalement et encore plus petites pour les V longitudinaux. La petitesse de tels signaux dans le modèle standard ouvre la porte à la recherche d'une nouvelle physique donnant lieu a de plus grands signaux dans ces modes.

C'est ce dont traite le troisième article.

N.B.: Dans les deux articles cités plus haut, on n'a pas mentionné la nuance concernant le signe entre A_T et $\overline{A}_T$ (qui dépend du type de PT) dans l'asymétrie de brisure T, $\mathcal{A}_T$. Mais cela n'affecte en rien nos résultats car on a calculé la bonne quantité (le vrai signal) donnée par l'éq. (49).

Dans le troisième article de cette thèse, on s'intéresse aux produits triples au-delà du modèle standard. On construit un lagrangien efficace pour les desintégrations non charmées de Λ_b, ayant pour sous-processus les modes $b \to s\bar{q}q$ et $b \to d\bar{q}q$. Ce lagrangien contient tous les termes possibles de dimension 6 de la nouvelle physique. Ensuite, on a calculé la contribution de ces termes aux asymétries T dues aux produits triples, toujours dans le cadre du modèle de factorisation. Cette méthode permet de voir quels opérateurs (donc quels modèles de la nouvelle physique) donnent de grandes asymétries T dans des modes où elles sont négligeables dans le modèle standard, ce qui permet de faire un tri parmi les opérateurs de la nouvelle physique. On trouve que la nouvelle physique change considérablement les résultats du modèle standard. Des asymétries négligeables dans ce dernier sont augmentées par les nouveaux opérateurs jusqu'à 50%.

Le quatrième article est un complément des trois autres. Il concerne l'étude d'asymétries de polarisation et "avant-arrière" ("*forward-backward*") dans le mode $b \to s\tau^+\tau^-$ avec les deux leptons polarisés (mentionnons que ce ne sont pas là des asymétries de brisure CP). Cela permet aussi de tester le modèle standard et de sonder la nouvelle physique. En effet, il y a un total de 31 asymétries dans ce mode, dont 9 sont mesurables dans le modèle standard (de l'ordre de 10% et plus). Cela permet de mesurer les coefficients de Wilson et la masse du quark b et de tester ainsi le modèle standard. Certaines des asymétries sont trouvées carrément nulles dans le modèle standard. Donc si elles sont mesurées ou encore si d'autres sont trouvées plus grandes que les valeurs prédites, ce sera un signal de présence d'une nouvelle physique. Un exemple intèressant est l'existence d'asymétries dues à des produits triples (précisons que ce sont de faux signaux de brisure CP car dûs seulement aux phases fortes) qui sont négligeables dans le modèle standard (de l'ordre de 10^{-2}). Leur observation constituerait un signal évident de la nouvelle physique.

CONTRIBUTION DE L'ÉTUDIANTE AUX DIFFÉRENTS ARTICLES

Ma contribution fut essentielle à chacun des articles.

- J'ai effectué tous les calculs du premier article.
- Dans le deuxième article, j'ai fait le calcul des amplitudes au carré dont les résultats sont représentés par les équations (6) et (42) et j'ai vérifié le reste des calculs.
- Dans le troisième article j'ai fais tous les calculs.
- J'ai effectué la plupart des calculs du quatrième article dont l'amplitude au carré et les différentes assymétries de polarisation.

CHAPITRE I: PREMIER ARTICLE

(Publié dans: Physical Review D 64, 116003 (2001).)

T-Violating Triple-Product Correlations in Hadronic b Decays

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Abstract

We study T-violating triple-product asymmetries in the quark-level decay $b \to su\bar{u}$ within the standard model (SM). We find that only two types of triple products are non-negligible. First, the asymmetry in $\vec{p}_u \cdot (\vec{s}_u \times \vec{s}_{\bar{u}})$ or $\vec{p}_{\bar{u}} \cdot (\vec{s}_u \times \vec{s}_{\bar{u}})$ can be as large as about 5%. It can be probed in $B \to V_1 V_2$ decays, where V_1 and V_2 are vector mesons. Second, the asymmetry in $\vec{s}_b \cdot (\vec{p}_u \times \vec{p}_s)$ can be in the range 1%–3%. One can search for this signal in decays such as $\Lambda_b \to \Lambda \pi^+ \pi^-$. All other triple-product asymmetries are expected to be small within the SM. This gives us new methods of searching for new physics.

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1 Introduction

These are exciting times for B physics. Measurements of $\sin 2\beta$ have been made [1], and provide the first hints of CP violation outside the kaon system. It is expected that further measurements of CP-violating rate asymmetries in the B system will be made before too long. And in the near future, data from HERA-B and hadron colliders will add to our knowledge of CP violation in the B system.

The purpose of all this activity is to test the standard model (SM) explanation of CP violation. In the SM, CP violation, which to date has been only seen in the kaon system, is due to the presence of a nonzero complex phase in the Cabibbo-Kobayashi-Maskawa (CKM) quark mixing matrix V. In this scenario, one expects large CP-violating effects in B decays, and the above experiments are searching for such signals.

The CP-violating signals which have been the most extensively studied are rate asymmetries in B decays [2]. Measurements of such asymmetries will allow one to cleanly probe the interior angles α, β and γ of the unitarity triangle, which will in turn provide important tests of the SM.

However, there is another class of CP-violating signals which has received relatively little attention: triple-product correlations [3]. In a given decay, it may be possible to measure the momenta and/or spins of the particles involved. From these one can construct triple products of the form $\vec{v}_1 \cdot (\vec{v}_2 \times \vec{v}_3)$, where each v_i is a spin or momentum. Such triple products are odd under time reversal (T) and hence, by the CPT theorem, are also potential signals of CP violation.

To establish the presence of a nonzero triple-product correlation, one constructs an asymmetry of the form

$$A_T \equiv \frac{\Gamma(\vec{v}_1 \cdot (\vec{v}_2 \times \vec{v}_3) > 0) - \Gamma(\vec{v}_1 \cdot (\vec{v}_2 \times \vec{v}_3) < 0)}{\Gamma(\vec{v}_1 \cdot (\vec{v}_2 \times \vec{v}_3) > 0) + \Gamma(\vec{v}_1 \cdot (\vec{v}_2 \times \vec{v}_3) < 0)} , \quad (1)$$

where Γ is the decay rate for the process in question. However, there is a complication: although the action of T changes the sign of a triple product, if a triple product changes sign, it is not necessarily due to the T transformation. This is because, in addition to reversing spins and momenta, the time reversal symmetry T also exchanges the initial and final states. Thus, in a given decay, a nonzero triple product is not necessarily a signal of T (and CP) violation. In particular, triple-product correlations can be faked by the presence of strong phases, even if there is no CP violation. That is, one typically finds that

$$A_T \propto \sin(\phi + \delta), \quad (2)$$

where ϕ is a weak, CP-violating phase and δ is a strong phase. From this we see that if $\delta \neq 0$, a triple-product correlation will appear, even in the absence of CP violation (i.e. if $\phi = 0$). In what follows, we refer to the triple-product asymmetries of Eq. (1) as *T-odd* effects.

Nevertheless, one can construct a *T-violating asymmetry*:

$$\mathcal{A}_T \equiv \frac{1}{2}(A_T - \bar{A}_T) , \quad (3)$$

where $\bar{A}_T$ is the T-odd asymmetry measured in the CP-conjugate decay process. This is a true T-violating signal in that it is nonzero only if $\phi \neq 0$ (i.e. if CP violation is present). Furthermore, unlike decay-rate asymmetries in direct CP violation, a nonzero $\mathcal{A}_T$ does not require the presence of a nonzero strong phase. Indeed:

$$\mathcal{A}_T \propto \sin\phi\cos\delta , \quad (4)$$

so that the signal is maximized when the strong phase is zero.

As with all CP-violating signals, (at least) two decay amplitudes are necessary to produce a triple-product correlation. Such correlations have been studied in semileptonic B decays [4]. However, since there is only a single amplitude in the SM, any such signal can occur only in the presence of new physics.

To our knowledge, the only study of triple products in the SM has been made by Valencia [5], who examined the decay $B \to V_1 V_2$, where V_1 and V_2 are vector mesons. He looked at triple products of the form $\vec{k} \cdot (\vec{\epsilon}_1 \times \vec{\epsilon}_2)$, where $\vec{\epsilon}_1$ and $\vec{\epsilon}_2$ are the polarizations of V_1 and V_2, respectively, and $\vec{k}$ is the momentum of one of the vector mesons. Since the calculation was done at the meson level, estimates of the various form factors were needed. The conclusion of this study was that, within the SM, one could expect a T-violating asymmetry at the level of several percent.

In this paper we re-examine the question of triple products in the SM using a complementary approach. In particular, we search for triple-product correlations at the quark level. The motivation is the following: if a significant triple-product correlation exists at the hadron level, it must also exist at the quark level. After all, given that QCD (which is responsible for hadronization) is CP-conserving, it is difficult to see how one can generate a large T-violating asymmetry at the hadron level if it is absent at the quark level.

Of course, the converse is not necessarily true: a large T-violating effect at the quark level might be "washed out" to some extent during hadronization, since the spins and momenta of the quarks may not correlate well with the spins and momenta of the hadrons. (The most obvious example of this is if spin-0 mesons are involved. In this case no information about the spins of the constituent quarks can be obtained.) Thus, a quark model is required to relate the asymmetries at the quark and meson levels. This suggests that the study of triple products at both the quark and meson levels may allow us to distinguish among the various quark models.

With this in mind, in this paper we examine the inclusive decay $b \to su\bar{u}$ within the SM. If there is a large $\vec{k} \cdot (\vec{\epsilon}_1 \times \vec{\epsilon}_2)$ triple product in $B \to V_1 V_2$, there should also be a large triple product at the quark level of the form $\vec{p} \cdot (\vec{s} \times \vec{s}')$, where $\vec{p}$ is the momentum of one of the quarks, and $\vec{s}$ and $\vec{s}'$ are the spins of two of the light quarks. And indeed, we find that the quark-level T-violating asymmetry due to the triple product $\vec{p}_u \cdot (\vec{s}_u \times \vec{s}_{\bar{u}})$ or $\vec{p}_{\bar{u}} \cdot (\vec{s}_u \times \vec{s}_{\bar{u}})$ is about 5%. This strongly supports Valencia's conclusion that the SM predicts a measurable T-violating asymmetry in $B \to V_1 V_2$.

However, we also find another significant T-violating signal in $b \to su\bar{u}$. It is due to the triple-product $\vec{s}_b \cdot (\vec{p}_u \times \vec{p}_s)$, which involves the b-quark spin and the momenta of the s and u quarks. In the SM, this signal turns out to be in the range of 1% to 3% of the total rate, which may be measurable. It might be observable in decays such as $\Lambda_b \to \Lambda \pi^+ \pi^-$.

Finally, it is also important to note which T-violating signals are *not* present. For example, we find that there are no significant T-violating asymmetries in the SM which involve the spin of the s-quark. Thus, should such an asymmetry be measured, it would be a clear sign of new physics.

In Sec. 2, we compute the triple products present in the decay $b \to su\bar{u}$, and estimate their sizes. We discuss possible hadron-level applications in Sec. 3. We conclude in Sec. 4.

2 Triple Products in $b \to su\bar{u}$

In the inclusive decay $b \to su\bar{u}$, the amplitude has two dominant contributions: the tree diagram (T) due to W-boson exchange and the loop-level strong penguin diagram (P). Furthermore, the penguin amplitude contains two dominant terms, P_1 and P_2 [6]. These various contributions are given by:

$$\begin{aligned} T &= \frac{4G_F}{\sqrt{2}} V_{ub}V_{us}^* \left[\bar{u}\gamma_\mu\gamma_L b\right]\left[\bar{s}\gamma^\mu\gamma_L v_u\right] e^{i\delta_t} , \\ P_1 &= -\frac{\alpha_s G_F}{\sqrt{2}\pi} F_1^c V_{cb}V_{cs}^* \left[\bar{s}t^\alpha\gamma_\mu\gamma_L b\right]\left[\bar{u}t_\alpha\gamma^\mu v_u\right] e^{i\delta_1} , \\ P_2 &= -\frac{\alpha_s G_F}{\sqrt{2}\pi} \left[\frac{-im_b}{q^2}F_2\right] V_{tb}V_{ts}^* \left[\bar{s}t^\alpha\sigma_{\mu\nu}q^\nu\gamma_R b\right]\left[\bar{u}t_\alpha\gamma^\mu v_u\right] e^{i\delta_2} . \end{aligned} \quad (5)$$

In the above, $\gamma_{L(R)} = (1 \mp \gamma_5)/2$, the t^α are the Gell-Mann matrices, and the δ_i are the strong phases. In P_2, q is the momentum of the internal gluon. The factors F_1^c and F_2 are functions of (m_c^2/M_W^2) and (m_t^2/M_W^2), respectively, and take the values $F_1^c \simeq 5.0$ and $F_2 \simeq 0.2$ for $m_t = 160$ GeV [6]. P_1 and P_2 are often called the *chromoelectric dipole moment* term and *chromomagnetic dipole moment* term, respectively.

The next step is the calculation of the square of the decay amplitude. We have:

$$|\mathcal{M}|^2 = |T|^2 + |P_1|^2 + |P_2|^2 + 2Re\left(T^\dagger P_1\right) + 2Re\left(T^\dagger P_2\right) + 2Re\left(P_1^\dagger P_2\right) . \quad (6)$$

The dominant term here is $|P_1|^2$.

We find triple products in all of the interference terms above (i.e. the last three terms of $|\mathcal{M}|^2$). Before giving the specific forms of these triple products, we make the following general remarks:

- In the calculation, we neglect the masses of the light quarks s, u and $\bar{u}$, but we keep the spins (i.e. polarization four-vectors) of these particles (at least to begin with). It turns out that there are no triple products involving the polarization of the s quark. (In other words, such terms are suppressed by at least m_s/m_b.) In light of this, in our results below, we automatically sum over the s-quark spin states.

This is an interesting result: it suggests that if a triple product involving the s-quark polarization is observed experimentally, it is probably due to physics beyond the SM.

- Since the B meson has spin 0, triple products in $B \to V_1 V_2$ cannot involve the spin of the b-quark. If one sums over the spin of the b-quark, the only term which contains triple products is the $T - P_1$ interference term. We will therefore use this term to estimate the size of the T-violating asymmetry in $B \to V_1 V_2$ [5].

- If the spins of the u and $\bar{u}$ quarks cannot be measured, one can then take them to be unpolarized, i.e. we sum over their polarizations. In this case, only the $T - P_2$ and $P_1 - P_2$ interferences contain a triple product. This unique signal takes the form $\vec{s}_b \cdot (\vec{p}_u \times \vec{p}_s)$.

- In all interference terms, there are triple products which involve the three polarizations $\vec{s}_b$, $\vec{s}_u$ and $\vec{s}_{\bar{u}}$. Experimentally, such signals will be extremely difficult to measure, and so are of less interest than the others described here.

2.1 $T - P_1$ interference

Keeping explicit the spins of the b-, u- and $\bar{u}$-quarks, the T-odd piece of the $T - P_1$ interference term is

$$\left[\sum_{s\ spins} 2Re\left(T^\dagger P_1\right)\right]_{T-odd} = \frac{16\alpha_s G_F^2 F_1^c}{3\pi}\, Im\left[V_{cs}^* V_{cb} V_{us} V_{ub}^* e^{i(\delta_1-\delta_t)}\right]$$
$$\times \Bigg\{ 2(p_b \cdot s_u)\, \epsilon_{\mu\nu\rho\xi}\, p_b^\mu p_u^\nu p_{\bar{u}}^\rho s_{\bar{u}}^\xi - 2(p_b \cdot p_u)\, \epsilon_{\mu\nu\rho\xi}\, p_b^\mu s_u^\nu p_{\bar{u}}^\rho s_{\bar{u}}^\xi + m_b^2\, \epsilon_{\mu\nu\rho\xi}\, p_u^\mu s_u^\nu p_{\bar{u}}^\rho s_{\bar{u}}^\xi$$
$$+ m_b \Big[(s_b \cdot s_u)\, \epsilon_{\mu\nu\rho\xi}\, p_s^\mu p_u^\nu p_{\bar{u}}^\rho s_{\bar{u}}^\xi - (s_b \cdot p_u)\, \epsilon_{\mu\nu\rho\xi}\, p_s^\mu p_{\bar{u}}^\nu s_{\bar{u}}^\rho s_u^\xi$$
$$- (p_s \cdot p_{\bar{u}})\, \epsilon_{\mu\nu\rho\xi}\, p_u^\mu s_b^\nu s_u^\rho s_{\bar{u}}^\xi + (p_s \cdot s_{\bar{u}})\, \epsilon_{\mu\nu\rho\xi}\, p_u^\mu p_{\bar{u}}^\nu s_b^\rho s_u^\xi \Big] \Bigg\} \quad (7)$$

Here, p_i is the 4-momentum of the i-quark and s_i is its polarization four-vector. Triple products[3] are found in the terms $\epsilon_{\mu\nu\rho\xi} v_1^\mu v_2^\nu v_3^\rho v_4^\xi$.

In the above expression, we see that there are two categories of triple products: those which involve s_b, the b-quark polarization, and those which do not. Those terms which include s_b [the last four terms in Eq. (7)] also include the polarizations of the u- and $\bar{u}$-quarks (s_u and $s_{\bar{u}}$). Since all three spins must be measured, these triple products will be extremely difficult to observe experimentally. Because of this, it is the first three terms of Eq. (7) which most interest us, and we therefore isolate them by averaging over s_b.

Of course, as written, the terms $\epsilon_{\mu\nu\rho\xi} v_1^\mu v_2^\nu v_3^\rho v_4^\xi$ involve only four-vectors, and therefore do not look like triple products. In order to identify the triple products implicit in these terms, we have to choose a particular frame of reference. The most natural choice is the rest frame of the b-quark, in which case Eq. (7) then takes the form

$$\left[\frac{1}{2}\sum_{b,s\ spins} 2Re\left(T^\dagger P_1\right)\right]_{T-odd} = \frac{16\alpha_s G_F^2 F_1^c}{3\pi} Im\left[V_{cs}^* V_{cb} V_{us} V_{ub}^* e^{i\Delta_{1t}}\right] m_b^2$$
$$\times \left\{ s_u^0 \vec{p}_u \cdot (\vec{p}_{\bar{u}} \times \vec{s}_{\bar{u}}) + E_u \vec{p}_{\bar{u}} \cdot (\vec{s}_u \times \vec{s}_{\bar{u}}) + s_{\bar{u}}^0 \vec{p}_u \cdot (\vec{p}_{\bar{u}} \times \vec{s}_u) + E_{\bar{u}} \vec{p}_u \cdot (\vec{s}_u \times \vec{s}_{\bar{u}}) \right\} \quad (8)$$

where $\Delta_{1t} \equiv \delta_1 - \delta_t$. We therefore see that there are, in fact, four distinct triple products in the $T - P_1$ interference term. These triple products depend on the polarization four-vectors of the u- and $\bar{u}$-quarks, whose most general form is [7]

$$s_i^\mu = \left(\frac{\vec{n}_i \cdot \vec{p}_i}{m_i} \ , \ \vec{n}_i + \frac{\vec{n}_i \cdot \vec{p}_i}{m_i(E_i + m_i)} \vec{p}_i \right) , \quad (9)$$

for $i = u, \bar{u}$. In the above, $\vec{n}_i$ is the polarization vector of the i-quark in its rest frame, and satisfies $|\vec{n}_i| = 1$.

[3]Note that, due to the identity $g_{\alpha\beta}\,\epsilon_{\mu\nu\rho\xi} - g_{\alpha\mu}\,\epsilon_{\beta\nu\rho\xi} - g_{\alpha\nu}\,\epsilon_{\mu\beta\rho\xi} - g_{\alpha\rho}\,\epsilon_{\mu\nu\beta\xi} - g_{\alpha\xi}\,\epsilon_{\mu\nu\rho\beta} = 0$, not all terms of the form $v_1 \cdot v_2\, \epsilon_{\mu\nu\rho\xi} v_3^\mu v_4^\nu v_5^\rho v_6^\xi$ are necessarily independent.

In order to compute the size of these triple-product asymmetries, in addition to integrating over phase space, we also need estimates of the sizes of the weak and strong phases. In the Wolfenstein parametrization [8], we can write the T-odd combination of CKM and strong phases as

$$Im\left[V_{cs}^* V_{cb} V_{us} V_{ub}^* e^{i\Delta_{1t}}\right] = A^2\lambda^6\left[\eta\cos\Delta_{1t} + \rho\sin\Delta_{1t}\right]. \tag{10}$$

CP violation in the CKM matrix is parametrized by the parameter η. As discussed in the introduction, nonzero strong phases can fake a T-violating signal. The term $\rho\sin\Delta_{1t}$ in the above expression is an example of such a fake signal. However, by forming a true T-violating asymmetry $\mathcal{A}_T$ [Eq. (3)], one can eliminate this fake signal. In this case $\mathcal{A}_T \propto \eta\cos\Delta_{1t}$.

At present, η is constrained to lie in the range $0.2 \leq \eta \leq 0.5$ [9]. As for the strong phase, the tree-level phase δ_t is usually assumed to be small: the logic is that, roughly speaking, the quarks will hadronize before having time to exchange gluons. On the other hand, for the $b \to su\bar{u}$ penguin amplitude, it is often assumed that strong phases come from the absorptive part of the penguin contribution [10]. Since P_1 involves an internal c-quark, it is possible that $\delta_1 \neq 0$, which of course implies that $\Delta_{1t} \neq 0$. Even so, for simplicity, in our calculation we assume that Δ_{1t} is small enough that $\cos\Delta_{1t} \simeq 1$ is a good approximation. However, the reader should be aware that the asymmetries may be reduced should this strong phase be large. (Note that the T-violating signal is maximal when $\cos\Delta_{1t} = 1$. For comparison, direct CP-violating rate asymmetries require the strong phase to be nonzero.)

We have performed the phase-space integration using the computer program RAMBO. In calculating the four T-violating asymmetries [see Eq. (3)] $\mathcal{A}_T^1$, $\mathcal{A}_T^2$, $\mathcal{A}_T^3$, and $\mathcal{A}_T^4$, which correspond respectively to the four triple products of Eq. (8): $s_u^0\,\vec{p}_u\cdot(\vec{p}_{\bar{u}}\times\vec{s}_{\bar{u}})$, $E_u\,\vec{p}_{\bar{u}}\cdot(\vec{s}_u\times\vec{s}_{\bar{u}})$, $s_{\bar{u}}^0\,\vec{p}_u\cdot(\vec{p}_{\bar{u}}\times\vec{s}_u)$, and $E_{\bar{u}}\,\vec{p}_u\cdot(\vec{s}_u\times\vec{s}_{\bar{u}})$, we have averaged over all directions of $\vec{s}_u$ and $\vec{s}_{\bar{u}}$. Our results are as follows.

$$\langle\mathcal{A}_T^{1,3}\rangle \simeq 0\ ,$$

$$2.3\% \lesssim \langle \mathcal{A}_T^{2,4} \rangle \lesssim 5.6\% \ , \tag{11}$$

where the range of $\mathcal{A}_T^{2,4}$ is due to the presently-allowed range for η.

Note that the triple product in $B \to V_1 V_2$ discussed by Valencia [5] is of the form $\vec{k} \cdot (\vec{\epsilon}_1 \times \vec{\epsilon}_2)$. In Eq. (8), it is the terms $\vec{p}_{\bar{u}} \cdot (\vec{s}_u \times \vec{s}_{\bar{u}})$ and $\vec{p}_u \cdot (\vec{s}_u \times \vec{s}_{\bar{u}})$ which could potentially give such a triple-product signal. We have found that the asymmetries $\mathcal{A}_T^2$ and $\mathcal{A}_T^4$, which correspond to these triple products, can be reasonably big ($\lesssim 5\%$). This is consistent with the results found by Valencia at the meson level, and suggests that the SM does indeed predict a measurable T-violating asymmetry in $B \to V_1 V_2$ decays.

Finally, for comparison, consider the decay-rate asymmetry, calculated by Hou for the same process [11]:

$$a_{CP}(b \to s u \bar{u}) \simeq 1.4\% \tag{12}$$

We therefore see that one expects T-violating triple-product asymmetries in $b \to s u \bar{u}$ to be considerably larger than the decay rate asymmetry.

2.2 $P_1 - P_2$ interference

The T-odd piece of the $P_1 - P_2$ interference term is

$$\begin{aligned}
\left[\sum_{s\ spins} 2Re\left(P_1^\dagger P_2\right) \right]_{T-odd} &= \frac{4\alpha_s^2 G_F^2 F_1^c F_2 m_b}{3\pi^2 q^2} Im\left[V_{ts}^* V_{tb} V_{cs} V_{cb}^* e^{i(\delta_2 - \delta_1)} \right] \\
&\times \Bigg\{ \Big[p_b \cdot (p_u - p_{\bar{u}}) \left(1 - s_u \cdot s_{\bar{u}}\right) - (s_{\bar{u}} \cdot p_s)(s_u \cdot p_{\bar{u}}) \\
&+ (s_u \cdot p_s)(s_{\bar{u}} \cdot p_u) \Big] \epsilon_{\mu\nu\rho\xi}\, p_b^\mu s_b^\nu p_u^\rho p_s^\xi \\
&+ \left[(s_u \cdot p_{\bar{u}})(p_u \cdot p_s) - \frac{q^2}{2}(s_u \cdot p_s) \right] \epsilon_{\mu\nu\rho\xi}\, p_b^\mu s_b^\nu p_s^\rho s_{\bar{u}}^\xi \\
&+ \left[(s_{\bar{u}} \cdot p_u)(p_{\bar{u}} \cdot p_s) - \frac{q^2}{2}(s_{\bar{u}} \cdot p_s) \right] \epsilon_{\mu\nu\rho\xi}\, p_b^\mu s_b^\nu p_s^\rho s_u^\xi \Bigg\} \ .
\end{aligned} \tag{13}$$

Here, if we average over the b-quark spin states, there is no T-violating signal at all.

We note that most of the terms in Eq. (13) correspond to triple products in which three spins must be measured. As we have already discussed, such signals are very difficult to observe experimentally, and so do not interest us. There is one term, however, which does not involve three spins, and it can be isolated by summing over the u- and $\bar{u}$-quark spin states:

$$\left[\sum_{u,\bar{u},s\ spins} 2Re\left(P_1^\dagger P_2\right)\right]_{T-odd} = \frac{16\alpha_s^2 G_F^2 F_1^c F_2 m_b}{3\pi^2 q^2} Im\left[V_{ts}^* V_{tb} V_{cs} V_{cb}^* e^{i(\delta_2-\delta_1)}\right] \times p_b \cdot (p_u - p_{\bar{u}})\, \epsilon_{\mu\nu\rho\xi} p_b^\mu s_b^\nu p_u^\rho p_s^\xi \,. \quad (14)$$

In the rest frame of the b-quark, the triple product takes the form $m_b^2(E_u - E_{\bar{u}})\, \vec{s}_b \cdot (\vec{p}_u \times \vec{p}_s)$. Integrating over phase space with RAMBO, we find that the $P_1 - P_2$ T-violating asymmetry is $O(10^{-5})$, which is negligible.

2.3 $T - P_2$ interference

Like $P_1 - P_2$ interference, the $T - P_2$ interference term contains two types of triple produts: (i) those involving a single quark polarization, s_b, and (ii) those involving the three polarization four-vectors s_b, s_u and $s_{\bar{u}}$. As usual, we are not interested in triple products involving three spins, and so we can therefore sum over s_u and $s_{\bar{u}}$. The T-odd piece of the $T - P_2$ interference term is then given by

$$\left[\sum_{u,\bar{u},s\ spins} 2Re\left(T^\dagger P_2\right)\right]_{T-odd} = \frac{128\alpha_s G_F^2 F_2 m_b}{3\pi q^2} Im\left[V_{ts}^* V_{tb} V_{us} V_{ub}^* e^{i\Delta_{2t}}\right] \times p_s \cdot p_u\, \epsilon_{\mu\nu\rho\xi}\, p_b^\mu s_b^\nu p_u^\rho p_s^\xi \,, \quad (15)$$

where $\Delta_{2t} \equiv \delta_2 - \delta_t$.

As was the case for the $T - P_1$ interference term, the T-violating asymmetry $\mathcal{A}_T$ is proportional to $\eta \cos \Delta_{2t}$. And, as before, we expect the δ_t piece of the strong phase Δ_{2t} to be small. However, there is a difference here compared to the $T - P_1$ case: previously, the penguin amplitude P_1 involved an internal c-quark, and so it was possible that the strong phase δ_{1t}, which is related to the absorptive part of the amplitude, could be sizeable. Here, the triple product involves only the t-quark penguin contribution P_2, which is purely dispersive, and so leads to $\delta_2 = 0$. Thus, it is an excellent approximation to set $\Delta_{2t} \simeq 0$.

In the rest frame of the b-quark, the triple-product of Eq. (15) is $m_b \, p_s \cdot p_u \, \vec{s}_b \cdot (\vec{p}_u \times \vec{p}_s)$. Integrating over phase space using RAMBO, and using the allowed range for η, we find that the corresponding T-violating triple-product asymmetry $\mathcal{A}_T^b$ can be of the order of several percent:

$$1.3\% \lesssim \mathcal{A}_T^b \lesssim 3.2\% \; . \qquad (16)$$

This could conceivably be measured at a future experiment.

Furthermore, if it is found that this asymmetry is considerably larger than the above values, it is probably a signal of new physics. For example, in some models of new physics, the chromomagnetic dipole moment F_2 can be enhanced up to ten times its SM value [12]. This will clearly have an enormous affect on the above asymmetry.

3 Applications

In the previous section, in our study of the quark-level decay $b \to s u \bar{u}$ within the SM, we found two classes of triple products whose T-violating asymmetry is large. They are: (i) $E_u \, \vec{p}_{\bar{u}} \cdot (\vec{s}_u \times \vec{s}_{\bar{u}})$ and $E_{\bar{u}} \, \vec{p}_u \cdot (\vec{s}_u \times \vec{s}_{\bar{u}})$, and (ii) $m_b \, p_s \cdot p_u \, \vec{s}_b \cdot (\vec{p}_u \times \vec{p}_s)$. The next obvious question is then: how can one test these results?

The ideal way would be to make triple-product measurements *inclusively*. If this were possible, then it would be straightforward to compare the experimental values with the theoretical predictions. However, this may not be experimentally feasible, in which case we must turn to exclusive B decays.

The first class of triple-product asymmetries can be studied in $B \to V_1 V_2$ decays which are dominated by the quark-level process $b \to s u \bar{u}$. Examples of such decays include $\overline{B_d^0} \to \rho K^*$, $\overline{B_s^0} \to K^{*+} K^{*-}$, $B_c^- \to D^* K^{*-}$, etc. These have been examined by Valencia, and we refer the reader to Ref. [5] for details.

Turning to the second class of triple products, it is clear that we cannot use decays of B mesons to obtain these asymmetries: since the B-meson spin is zero, there is no way to measure the spin of the b-quark (which is the only spin contributing to the triple product). However, one possibility would be to use the Λ_b baryon, whose spin is largely that of the b quark. For example, we can consider the process $\Lambda_b \to \Lambda \pi^+ \pi^-$. The triple product $\vec{s}_b \cdot (\vec{p}_u \times \vec{p}_s)$ can be roughly equated to $\vec{s}_{\Lambda_b} \cdot (\vec{p}_{\pi^+} \times \vec{p}_\Lambda)$.

In all cases, since the underlying quark-level process is $b \to s u \bar{u}$, we expect the branching ratios for the B or Λ_b decays to be $O(10^{-5})$. However, for certain decays, there may be a dynamical suppression, and this may affect the expected triple-product asymmetry. For example, suppose that the $\pi^+\pi^-$ pair in $\Lambda_b \to \Lambda \pi^+ \pi^-$ comes mainly from a virtual ρ-meson. In this case, due to isospin conservation, the gluonic penguin will not contribute to this decay, which means that the branching ratio will be considerably smaller than $O(10^{-5})$ [13]. More importantly, the triple-product asymmetry will vanish, since we no longer have two interfering amplitudes. Should this occur, a better decay mode in which to look for the $\vec{s}_b \cdot (\vec{p}_u \times \vec{p}_s)$ triple product might be $\Lambda_b \to \Lambda K^+ K^-$, which should not suffer a similar suppression. The main lesson here is that it is important to search for triple products in a variety of decays.

For both types of triple-product asymmetries, it will be necessary to use a quark model to make the connection between the experimental results and the quark-level calculations. In particular, we will want to know how the spin and momentum of a hadron is related to the spin and momentum of the constituent quarks. Some of these relations are on relatively firm footing. For example, it is widely accepted that the spin of the Λ_b (Λ) is essentially equal to the spin of the internal b-quark (s-quark) [14]. On the other hand, different quark models predict different relationships between the spin of a vector meson and the spin of its internal quarks, and similarly for the momenta of mesons and quarks. Thus, the study of triple-product correlations may allow us to distinguish among the various quark models which have been proposed.

Finally, we note that certain quark-level triple products are predicted to be small in the SM. For example, triple products involving the spin of the s-quark are suppressed by powers of its mass. Hence, if a T-violating asymmetry due to a triple product involving the s-quark spin were found to be sizeable, this would probably indicate the presence of new physics. The decay $\Lambda_b \to \Lambda\pi^+\pi^-$, which was mentioned above, can be used to test this. The spin of the Λ is essentially equal to the s-quark spin, so any T-violating asymmetry involving the spin of the Λ, such as $\vec{s}_{\Lambda_b} \cdot (\vec{s}_\Lambda \times \vec{p}_\Lambda)$, should be tiny in the SM.

As another example, recall that we found that $T - P_1$ interference produced the triple products $s_u^0\, \vec{p}_u \cdot (\vec{p}_{\bar{u}} \times \vec{s}_{\bar{u}})$ and $s_{\bar{u}}^0\, \vec{p}_u \cdot (\vec{p}_{\bar{u}} \times \vec{s}_u)$. However, the corresponding T-violating asymmetries $\mathcal{A}_T^1$ and $\mathcal{A}_T^3$ turned out to be suppressed dynamically. Consider the decay of a B-meson to two vector mesons, $B \to V_1 V_2$, where the V_2 then subsequently decays to two mesons $\Phi_1\Phi_2$. Roughly speaking, one can relate $s_u^0\, \vec{p}_u \cdot (\vec{p}_{\bar{u}} \times \vec{s}_{\bar{u}})$ to $\epsilon_{V_1}^0\, \vec{p}_{V_1} \cdot (\vec{\epsilon}_{V_2} \times \vec{p}_{\Phi_1})$. Thus, the measurement of a nonzero value for this latter triple-product asymmetry would be a signal for new physics.

4 Conclusions

We have calculated the quark-level triple-product correlations in the decay $b \to su\bar{u}$ within the standard model. Although several such triple products are present, we find that only two types lead to sizeable T-violating asymmetries.

The first type includes $\vec{p}_u \cdot (\vec{s}_u \times \vec{s}_{\bar{u}})$ and $\vec{p}_{\bar{u}} \cdot (\vec{s}_u \times \vec{s}_{\bar{u}})$. We find that the corresponding T-violating asymmetries can be as large as about 5%. This triple product can be probed in $B \to V_1 V_2$ decays, where V_1 and V_2 are vector mesons [5].

The second type is $\vec{s}_b \cdot (\vec{p}_u \times \vec{p}_s)$, where $\vec{s}_b$ is the polarization of the b-quark, and $\vec{p}_u$ and $\vec{p}_s$ are the momenta of the u- and s-quark, respectively. We calculate that the T-violating asymmetry for this triple product is in the range 1%–3%, which may be measurable. There are several ways to try to search for this triple-product asymmetry. For example, one could study the decay $\Lambda_b \to \Lambda\pi^+\pi^-$, looking for a nonzero triple product $\vec{s}_{\Lambda_b} \cdot (\vec{p}_{\pi^+} \times \vec{p}_\Lambda)$.

The fact that we find only two large triple-product correlations has interesting consequences. If a triple product is tiny at the quark level, it is probably tiny at the hadron level as well. After all, the hadronization of quarks into hadrons is a strong-interaction process, and QCD is CP-conserving. It is therefore difficult to see how one can generate a large triple-product correlation at the hadron level, given that it is small at the quark level. Thus, from the point of view of looking for physics beyond the SM, it is important to identify those triple-product asymmetries which are expected to be small in the SM. If such asymmetries are found to be large, this is probably a signal of new physics. For example, we find that triple products involving the spin of the s-quark are suppressed by powers of its mass. Thus, if, for instance, a sizeable T-violating asymmetry of the form $\vec{s}_{\Lambda_b} \cdot (\vec{s}_\Lambda \times \vec{p}_\Lambda)$ were found in the decay $\Lambda_b \to \Lambda\pi^+\pi^-$, this would be compelling evidence for the presence of new physics, since the spin of the Λ is due largely to the s-quark spin.

5 Acknowledgments

W.B. would like to thank P. Depommier for helpful discussions, and G. Azuelos for help with RAMBO. D.L. thanks A. Datta, B. Kayser and A. Soni for useful conversations. This work was financially supported by NSERC of Canada.

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CHAPITRE II: DEUXIÈME ARTICLE

(Publié dans: Physics Letters B 538, 309-320 (2002).)

T-violating Triple-Product Correlations in Charmless Λ_b Decays

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Abstract

Using factorization, we compute, within the standard model, the T-violating triple-product correlations in the charmless decays $\Lambda_b \to F_1 F_2$, where F_1 is a light spin-$\frac{1}{2}$ baryon and F_2 is a pseudoscalar (P) or vector (V) meson. We find a large triple-product asymmetry of 18% for the decay $\Lambda_b \to pK^-$. However, for other classes of $\Lambda_b \to F_1 P$ decays, the asymmetry is found to be at most at the percent level. For $\Lambda_b \to F_1 V$ decays, we find that all triple-product asymmetries are small (at most $O(1\%)$) for a transversely-polarized V, and are even smaller for longitudinal polarization. Our estimates of the nonfactorizable contributions to these decays show them to be negligible, and we describe ways of testing this.

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Over the past two decades, there has been a great deal of theoretical work examining CP violation in the B system. Most of this work has focussed on the decays of B mesons. The main reason is that the indirect CP-violating asymmetries in B-meson decays can be used to extract the interior angles of the unitarity triangle (α, β and γ) with no hadronic uncertainty [1]. The knowledge of these angles will allow us to test the standard model (SM) explanation of CP violation. In order to make such measurements, the B-factories BaBar and Belle have been built. These machines produce copious numbers of B^0–$\bar{B}^0$ pairs, and have now provided the first definitive evidence for CP violation outside the kaon system: $\sin 2\beta = 0.78 \pm 0.08$ [2].

On the other hand, in the coming years machines will be built which are capable of producing large numbers of Λ_b baryons. These include hadron machines, such as the Tevatron, LHC, etc., as well as possibly a high-luminosity e^+e^- machine running at the Z pole. People have therefore started to examine the SM predictions for a variety of Λ_b decays. This is a worthwhile effort, since it is conceivable that certain types of new physics will be more easily detectable in Λ_b decays than in B decays. For example, Λ_b's can be used to probe observables which depend on the spin of the b-quark, whereas such observables will be unmeasurable in B-meson decays.

One class of observables which may involve, among other things, the b-quark spin is triple-product correlations. These take the form $\vec{v}_1 \cdot (\vec{v}_2 \times \vec{v}_3)$, where each v_i is a spin or momentum. These triple products are odd under time reversal (T) and hence, by the CPT theorem, also constitute potential signals of CP violation. By measuring a nonzero value of the asymmetry

$$A_T \equiv \frac{\Gamma(\vec{v}_1 \cdot (\vec{v}_2 \times \vec{v}_3) > 0) - \Gamma(\vec{v}_1 \cdot (\vec{v}_2 \times \vec{v}_3) < 0)}{\Gamma(\vec{v}_1 \cdot (\vec{v}_2 \times \vec{v}_3) > 0) + \Gamma(\vec{v}_1 \cdot (\vec{v}_2 \times \vec{v}_3) < 0)} , \tag{1}$$

where Γ is the decay rate for the process in question, one can establish the presence of a nonzero triple-product correlation. Note that there is a well-known technical complication: strong phases can produce a nonzero value of A_T, even if there is no CP violation (i.e. if the weak phases are zero). Thus, strictly speaking, the asymmetry A_T is not in fact a T-violating effect. Nevertheless, one can still obtain a true T-violating signal by measuring a nonzero value of

$$\mathcal{A}_T \equiv \frac{1}{2}(A_T - \bar{A}_T) \ , \tag{2}$$

where $\bar{A}_T$ is the T-odd asymmetry measured in the CP-conjugate decay process.

Recently, T-violating triple-product correlations were calculated for the inclusive quark-level decay $b \to s\bar{u}u$ [3]. In that calculation, all final-state masses were neglected. Ignoring triple products which involve three spins, only two non-negligible triple-product asymmetries were found. They are: (i) $\vec{p}_u \cdot (\vec{s}_u \times \vec{s}_{\bar{u}})$ or $\vec{p}_{\bar{u}} \cdot (\vec{s}_u \times \vec{s}_{\bar{u}})$, and (ii) $\vec{s}_b \cdot (\vec{p}_u \times \vec{p}_s)$. While the former triple product can be probed in $B \to V_1 V_2$ decays, where V_1 and V_2 are vector mesons, the latter can only be measured in Λ_b decays, since the spin of the b quark is involved.

In this paper we study, within the SM, the triple products in charmless two-body Λ_b decays which are generated by the quark-level transitions $b \to s\bar{u}u$ or $b \to d\bar{u}u$. These decays are of the type $\Lambda_b \to F_1 F_2$, where F_1 is a light spin-$\frac{1}{2}$ baryon, such as p, Λ, etc., and F_2 is a pseudoscalar (P) or vector (V) meson. Such T-violating triple-product correlations, along with other P-violating asymmetries, have been studied for hyperon decays [4], but relatively little work has been done to study CP violation in Λ_b decays.

The decays $\Lambda_b \to F_1 P$ are similar to hyperon decays. As we will see, there is a triple-product correlation in such decays of the form $\vec{s}_{\Lambda_b} \cdot (\vec{s}_{F_1} \times \vec{p})$, where $\vec{s}_{\Lambda_b}$ and $\vec{s}_{F_1}$ are the polarizations of the Λ_b and F_1, respectively, and $\vec{p}$ is the momentum of one of the final-state particles in the rest frame of the Λ_b. On the other hand, Λ_b decays can also include a vector meson in the final state, which is not kinematically accesible for hyperon decays. The decay $\Lambda_b \to F_1 V$

can give rise to a variety of triple-product correlations involving the spin of the Λ_b and/or V.

Many of these triple products involve the spin of the Λ_b. Perhaps the easiest way to obtain this quantity is to produce the Λ_b baryons in the decay of an on-shell Z boson. This is because, in the decay $Z \to b\bar{b}$, the b-quarks have a large average longitudinal polarization of about -94%. According to heavy-quark effective theory, this polarization is retained when a b-quark hadronizes into a Λ_b, and recent measurements of the average longitudinal polarization of b-flavored baryons produced in Z^0 decays (measured through their decay to $\Lambda_c \ell \nu_\ell X$) are consistent with this conclusion [5]. Thus, the so-called GigaZ option (2×10^9 Z bosons per year [6, 7]) of a high-luminosity e^+e^- collider running at the Z peak would be a particularly good environment for measuring triple-product correlations in Λ_b decays. However, even if the spin of the Λ_b cannot be measured at a given machine, some of the triple-product correlations in $\Lambda_b \to F_1 V$ do not involve the polarization of the initial state. Thus, triple products can be measured at a variety of facilities in which a large number of Λ_b baryons is produced.

We begin our analysis by studying the nonleptonic decay $\Lambda_b \to F_1 P$. The general form for this amplitude can be written as

$$\mathcal{M}_P = A(\Lambda_b \to F_1 P) = i\bar{u}_{F_1}(a + b\gamma_5)u_{\Lambda_b} \, . \tag{3}$$

In order to make contact with the conventional notation for hyperon decay, we note that, in the rest frame of the parent baryon, the decay amplitude reduces to

$$A(\Lambda_b \to F_1 P) = i\chi_{F_1}(S + P\vec{\sigma}\cdot\hat{p})\chi_{\Lambda_b} \, , \tag{4}$$

where $\hat{p}$ is the unit vector along the direction of the daughter baryon momentum, and $S = \sqrt{2m_{\Lambda_b}(E_{F_1} + m_{F_1})}a$ and $P = -\sqrt{2m_{\Lambda_b}(E_{F_1} - m_{F_1})}b$, where E_{F_1} and m_{F_1} are, respectively, the energy and mass of the final-state baryon

F_1. The decay rate and the various asymmetries are given by

$$\Gamma = \frac{\vec{p}}{8\pi {m_{\Lambda_b}}^2}(|S|^2 + |P|^2) ,$$
$$\alpha = \frac{2\,\mathrm{Re}(S^*P)}{|S|^2+|P|^2} , \quad \beta = \frac{2\,\mathrm{Im}(S^*P)}{|S|^2+|P|^2} , \quad \gamma = \frac{|S|^2-|P|^2}{|S|^2+|P|^2} . \qquad (5)$$

(Note: above, the quantities α, β and γ should not be confused with the CP phases of the unitarity triangle, which have the same symbols.)

The calculation of $|\mathcal{M}_P|^2$ in Eq. (3) yields

$$\begin{aligned} |\mathcal{M}_P|^2 &= (|a|^2 - |b|^2)\,(m_{F_1} m_{\Lambda_b} + p_{F_1}\cdot s_{\Lambda_b}\, p_{\Lambda_b}\cdot s_{F_1} - p_{F_1}\cdot p_{\Lambda_b}\, s_{F_1}\cdot s_{\Lambda_b}) \\ &\quad +(|a|^2 + |b|^2)\,(p_{F_1}\cdot p_{\Lambda_b} - m_{F_1} m_{\Lambda_b} s_{F_1}\cdot s_{\Lambda_b}) \\ &\quad +2\,\mathrm{Re}(ab^*)\,(m_{\Lambda_b} p_{F_1}\cdot s_{\Lambda_b} - m_{F_1} p_{\Lambda_b}\cdot s_{F_1}) \\ &\quad +2\,\mathrm{Im}(ab^*)\epsilon_{\mu\nu\rho\sigma} p^\mu_{F_1} s^\nu_{F_1} p^\rho_{\Lambda_b} s^\sigma_{\Lambda_b} . \end{aligned} \qquad (6)$$

It is the last term above which gives a triple-product correlation. (It corresponds to β in Eq. (5).) In the rest frame of the Λ_b, it takes the form $\vec{p}_{F_1}\cdot(\vec{s}_{F_1}\times\vec{s}_{\Lambda_b})$.

In order to estimate the size of this triple product, we will use factorization to calculate $\mathrm{Im}(ab^*)$ at the hadron level. The starting point is the SM effective hamiltonian for charmless hadronic B decays [8]:

$$H^q_{eff} = \frac{G_F}{\sqrt{2}}[V_{ub}V^*_{uq}(c_1 O^q_1 + c_2 O^q_2) - \sum_{i=3}^{10} V_{tb}V^*_{tq} c^t_i O^q_i] + h.c., \qquad (7)$$

where

$$\begin{aligned} O^q_1 &= \bar{q}_\alpha \gamma_\mu L u_\beta\, \bar{u}_\beta \gamma^\mu L b_\alpha \quad , \quad O^q_2 = \bar{q}\gamma_\mu L u\, \bar{u}\gamma^\mu L b , \\ O^q_{3(5)} &= \bar{q}\gamma_\mu L b \sum_{q'} \bar{q}'\gamma^\mu L(R) q' \quad , \quad O^q_{4(6)} = \bar{q}_\alpha \gamma_\mu L b_\beta \sum_{q'} \bar{q}'_\beta \gamma^\mu L(R) q'_\alpha , \\ O^q_{7(9)} &= \frac{3}{2}\bar{q}\gamma_\mu L b \sum_{q'} e_{q'} \bar{q}'\gamma^\mu R(L) q' \quad , \quad O^q_{8(10)} = \frac{3}{2}\bar{q}_\alpha\gamma_\mu L b_\beta \sum_{q'} e_{q'} \bar{q}'_\beta \gamma^\mu R(L) q'_\alpha . \end{aligned} \qquad (8)$$

In the above, q can be either a d or an s quark, depending on whether the decay is a $\Delta S = 0$ or a $\Delta S = -1$ process, $q' = d$, u or s, with $e_{q'}$ the corresponding electric charge, and $R(L) = 1 \pm \gamma_5$. The values of the Wilson coefficients c_i evaluated at the scale $\mu = m_b = 5$ GeV, for $m_t = 176$ GeV and $\alpha_s(m_Z) = 0.117$, are [9]:

$$\begin{aligned}
c_1 = -0.324 \quad &, \quad c_2 = 1.151 \ , \\
c_3^t = 0.017 \ , \quad c_4^t = -0.037 \quad &, \quad c_5^t = 0.010 \ , \quad c_6^t = -0.045 \ , \\
c_7^t = -1.24 \times 10^{-5} \quad &, \quad c_8^t = 3.77 \times 10^{-4} \ , \\
c_9^t = -0.010 \quad &, \quad c_{10}^t = 2.06 \times 10^{-3} \ .
\end{aligned} \tag{9}$$

In our analysis we will also consider the gluonic dipole operator in which the gluon splits into two quarks, giving the effective operator

$$H_{11} = i \frac{G_F}{\sqrt{2}} \frac{\alpha_s(\mu)}{2\pi k^2} m_b(\mu)\, c_{11}\, V_{tb} V_{ts}^* \, \bar{s}(p_s) \sigma_{\mu\nu} R T^a b(p_b) \bar{q}(p_2) \gamma^\mu T^a q(p_1) k^\nu \ , \tag{10}$$

where $k = p_b - ps$ and $c_{11} = 0.2$ [10]. It is often useful to write this in the Fierz-transformed form

$$\begin{aligned}
H_{11} &= -\frac{G_F}{\sqrt{2}} \frac{\alpha_s(\mu)}{16\pi} \frac{m_b^2(\mu)}{k^2} c_{11} \frac{N_c^2 - 1}{N_c^2} V_{tb} V_{ts}^* \\
&\quad \times \left[\delta_{\alpha\beta} \delta_{\alpha'\beta'} - \frac{2N_c}{N_c^2 - 1} T^a_{\alpha\beta} T^a_{\alpha'\beta'} \right] \sum_i T_i \ ,
\end{aligned} \tag{11}$$

where

$$\begin{aligned}
T_1 &= 2 \bar{s}_\alpha \gamma_\mu L q_\beta \bar{q}_{\alpha'} \gamma^\mu L b_{\beta'} - 4 \bar{s}_\alpha R q_\beta \bar{q}_{\alpha'} L b_{\beta'} \ , \\
T_2 &= 2 \frac{m_s}{m_b} \bar{s}_\alpha \gamma_\mu R q_\beta \bar{q}_{\alpha'} \gamma^\mu R b_{\beta'} - 4 \frac{m_s}{m_b} \bar{s}_\alpha L q_\beta \bar{q}_{\alpha'} R b_{\beta'} \ , \\
T_3 &= \frac{(p_b + p_s)_\mu}{m_b} [\bar{s}_\alpha \gamma^\mu L q_\beta \bar{q}_{\alpha'} R b_{\beta'} + \bar{s}_\alpha R q_\beta \bar{q}_{\alpha'} \gamma^\mu R b_{\beta'}] \ , \\
T_4 &= \frac{(p_b + p_s)_\mu}{m_b} [i \bar{s}_\alpha \sigma^{\mu\nu} R q_\beta \bar{q}_{\alpha'} \gamma_\nu R b_{\beta'} - i \bar{s}_\alpha \gamma_\nu L q_\beta \bar{q}_{\alpha'} \sigma^{\mu\nu} R b_{\beta'}] \ ,
\end{aligned} \tag{12}$$

in which we have defined $\sigma_{\mu\nu} = \frac{i}{2}[\gamma_\mu\gamma_\nu - \gamma_\nu\gamma_\mu]$.

We now apply the effective hamiltonian to specific exclusive Λ_b decays. We will focus on those processes for which factorization is expected to be a good approximation, namely colour-allowed decays. We begin with $\Lambda_b \to pK^-$, which is a $b \to s\bar{u}u$ transition. Factorization allows us to write

$$A(\Lambda_b \to pK^-) = \sum_{O,O'} \left\langle K^- \right| O \left|0\right\rangle \left\langle p\right| O' \left|\Lambda_b\right\rangle . \qquad (13)$$

It is straightforward to show that the operators in H^s_{eff} and H_{11} lead to two classes of terms in the decay amplitude: (a) $\langle K^-| \bar{s}\gamma^\mu(1 \pm \gamma_5)u |0\rangle \langle p| \bar{u}\gamma_\mu(1 \pm \gamma_5)b |\Lambda_b\rangle$, and (b) $\langle K^-| \bar{s}(1 \pm \gamma_5)u |0\rangle \langle p| \bar{u}(1 \pm \gamma_5)b |\Lambda_b\rangle$. For the first of these, we define the pseudoscalar decay constant f_K as

$$if_K q^\mu = \langle K| \bar{s}\gamma^\mu(1 - \gamma_5)u |0\rangle \ , \qquad (14)$$

where $q^\mu \equiv p^\mu_{\Lambda_b} - p^\mu_p = p^\mu_K$ is the four-momentum transfer. For the second, one can show that

$$\begin{aligned} \left\langle K^- \right| \bar{s}(1 \pm \gamma_5)u \left|0\right\rangle &= \mp\frac{f_K m_K^2}{m_s + m_u}, \\ \langle p| \bar{u}(1 \pm \gamma_5)b |\Lambda_b\rangle &= \frac{q^\mu}{m_b} \langle p| \bar{u}\gamma_\mu(1 \mp \gamma_5)b |\Lambda_b\rangle \ . \end{aligned} \qquad (15)$$

(In the second matrix element, we have neglected m_u compared to m_b.) Thus, factorization leads to the following form for the $\Lambda_b \to pK^-$ amplitude:

$$\begin{aligned} A(\Lambda_b \to pK^-) &= if_K q^\mu \langle p| \bar{u}\gamma_\mu(1 - \gamma_5)b |\Lambda_b\rangle X_K \\ &\quad + if_K q^\mu \langle p| \bar{u}\gamma_\mu(1 + \gamma_5)b |\Lambda_b\rangle Y_K \ . \end{aligned} \qquad (16)$$

Like any CP-violating observable, a nonzero triple product can arise only if there are two interfering amplitudes. This will occur only if both X_K and Y_K are nonzero. Since all the operators O_1–O_{10} involve a left-handed b-quark, it is clear that $X_K \neq 0$ in the SM. Furthermore, though it is less obvious, one can also have $Y_K \neq 0$. Consider, for example, the operator O_6 of Eq. (9). After performing Fierz transformations, this can be written as

$$O_6 \sim \bar{s}(1+\gamma_5)u\,\bar{u}(1-\gamma_5)b\ . \tag{17}$$

However, according to Eq. (15), $\langle p|\,\bar{u}(1-\gamma_5)b\,|\Lambda_b\rangle$ can be related to $\langle p|\,\bar{u}\gamma_\mu(1+\gamma_5)b\,|\Lambda_b\rangle$. Thus, Y_K receives contributions from operators such as O_6. We find

$$\begin{aligned} X_K &= \frac{G_F}{\sqrt{2}}\left[V_{ub}V_{us}^* a_2 - \sum_{q=u,c,t} V_{qb}V_{qs}^*(a_4^q + a_{10}^q) - V_{tb}V_{ts}^* a_d\left(1+\frac{2E_K}{m_b}\right)\right] , \\ Y_K &= -\frac{G_F}{\sqrt{2}}\left[\sum_{q=u,c,t} V_{qb}V_{qs}^*(a_6^q + a_8^q) + \frac{5}{4}V_{tb}V_{ts}^* a_d\right]\chi_K\ , \end{aligned} \tag{18}$$

with

$$\chi_K = \frac{2m_K^2}{(m_s+m_u)m_b}\ , \qquad a_d = \frac{\alpha_s(\mu)}{16\pi}\langle\frac{m_b^2(\mu)}{k^2}\rangle c_{11}\frac{N_c^2-1}{N_c^2}\ . \tag{19}$$

In the above, we have defined $a_i^q = c_i^q + \frac{c_{i+1}^q}{N_c}$ for i odd and $a_i^q = c_i^q + \frac{c_{i-1}^q}{N_c}$ for i even. We estimate the average gluon momentum in the dipole operator to be $\langle m_b^2/k^2\rangle = \int \phi_K(x) m_b^2/k^2 dx$, where the gluon momentum in the heavy-quark limit is $k^2 = m_b^2(1-x)$ and ϕ_K is the kaon light-cone distribution. Choosing the asymptotic form $\phi_K = 6x(1-x)$, we find $\langle m_b^2/k^2\rangle = 3$, which leads to $a_d = 0.0021$.

Now, the vector and axial-vector matrix elements between the Λ_b and p can be written in the general form

$$\begin{aligned} \langle p|\,\bar{u}\gamma^\mu b\,|\Lambda_b\rangle &= \bar{u}_p\left[f_1\gamma^\mu + i\frac{f_2}{m_{\Lambda_b}}\sigma^{\mu\nu}q_\nu + \frac{f_3}{m_{\Lambda_b}}q^\mu\right]u_{\Lambda_b} \\ \langle p|\,\bar{u}\gamma^\mu\gamma_5 b\,|\Lambda_b\rangle &= \bar{u}_p\left[g_1\gamma^\mu + i\frac{g_2}{m_{\Lambda_b}}\sigma^{\mu\nu}q_\nu + \frac{g_3}{m_{\Lambda_b}}q^\mu\right]\gamma_5 u_{\Lambda_b}\ , \end{aligned} \tag{20}$$

where the f_i and g_i are Lorentz-invariant form factors. Heavy-quark symmetry imposes constraints on these form factors. A systematic expansion of these form factors, including $1/m_b$ corrections, has been calculated [11]: in the $m_b \to \infty$ limit, one obtains the relations

$$f_1 = g_1 \ , \qquad f_2 = g_2 = f_3 = g_3 \ . \tag{21}$$

Using the above expressions, we find that the parameters a and b of Eq. (3) can be written as

$$\begin{aligned} a_K &= f_K(X_K + Y_K)\left[(m_{\Lambda_b} - m_p)f_1 + f_3\frac{m_K^2}{m_{\Lambda_b}}\right] \ , \\ b_K &= f_K(X_K - Y_K)\left[(m_{\Lambda_b} + m_p)g_1 - g_3\frac{m_K^2}{m_{\Lambda_b}}\right] \ . \end{aligned} \tag{22}$$

According to Eq. 6, the triple product in $\Lambda_b \to pK^-$ is proportional to $\mathrm{Im}(a_K b_K^*)$, which is in turn proportional to $\mathrm{Im}(X_K Y_K^*)$. Since X_K and Y_K are both nonzero, and have different weak phases [Eq. (18)], we expect a nonzero triple-product asymmetry in $\Lambda_b \to pK^-$ of the form $\vec{p}_p \cdot (\vec{s}_p \times \vec{s}_{\Lambda_b})$. At first sight, this appears to contradict the results of Ref. [3], since no triple products involving two spins were found in the quark-level decay $b \to s\bar{u}u$. However, note that Y_K is proportional to χ_K, which is formally suppressed by $1/m_b$. Thus, in the limit $m_b \to \infty$, one has $Y_K = 0$, so that the triple-product correlation will vanish. This agrees with the conclusions of Ref. [3], which neglects the masses of the final-state quarks (i.e. the limit $m_b \to \infty$ is implicitly assumed).

However, the key point is that, for finite m_b, χ_K is not small because of the presence of the chiral enhancement term $m_K^2/(m_s + m_u)$. In fact, for $m_s = 100$ MeV and $m_b = 5$ GeV, $\chi_K \sim 1$, and hence is clearly non-negligible. The triple-product asymmetry of $\vec{p}_p \cdot (\vec{s}_p \times \vec{s}_{\Lambda_b})$ may therefore be sizeable. Note that this triple product requires the measurement of both the Λ_b and the p polarizations. If the measurement of the proton polarization is not possible, one can instead consider a final state with an excited nucleon, such as $\Lambda_b \to N(1440)K^-$. In this case the polarization of the $N(1440)$ can be determined from its decay products. (Alternatively, one can consider the decay $\Xi_b \to \Sigma^+ K^-$, where Ξ_b has quark content *bus*.)

Note also that in Eq. (18) we have included the up- and charm-quark penguin pieces, proportional to $V_{ub}V_{us}^*$ and $V_{cb}V_{cs}^*$ respectively. These are generated by rescattering of the tree-level operators in the effective Hamiltonian in Eq. (7). As we will see, the contributions from these rescattering terms are very important. The coefficients associated with these terms are given by

$$c_{3,5}^i = -c_{4,6}^i/N_c = P_s^i/N_c \quad , \quad c_{7,9}^i = P_e^i \, , \quad c_{8,10}^i = 0 \, , \quad i = u, c \, , \qquad (23)$$

where N_c is the number of colours. The leading contributions to $P_{s,e}^i$ are given by $P_s^i = (\frac{\alpha_s}{8\pi})c_2(\frac{10}{9} + G(m_i, \mu, q^2))$ and $P_e^i = (\frac{\alpha_{em}}{9\pi})(N_c c_1 + c_2)(\frac{10}{9} + G(m_i, \mu, q^2))$, in which the function $G(m, \mu, q^2)$ takes the form

$$G(m, \mu, q^2) = 4 \int_0^1 x(1-x) \ln \frac{m^2 - x(1-x)q^2}{\mu^2} \, dx \, , \qquad (24)$$

where q is the momentum carried by the virtual gluon in the penguin diagram. Of course, we are really interested in the matrix elements of the various operators for the decay $\Lambda_b \to pK^-$, and so the coefficients in Eq. (23) should be understood to be

$$\bar{c}_i^{u,c} = \frac{\langle pK^- | \, c_i^{u,c}(q^2) O_i \, | \Lambda_b \rangle}{\langle pK^- | \, O_i \, | \Lambda_b \rangle} \, . \qquad (25)$$

We will henceforth drop the distinction between $\bar{c}_i^{u,c}$ and $c_i^{u,c}$, with the understanding that it is the $\bar{c}_i^{u,c}$ which appear in the amplitude.

The analysis of other colour-allowed Λ_b decays follows straightforwardly from that for $\Lambda_b \to pK^-$. For example, consider $\Lambda_b \to p\pi^-$, which is generated by the quark-level decay $b \to d\bar{u}u$. The amplitude for $\Lambda_b \to p\pi^-$ is given by Eq. (3), with

$$a_\pi \;\; = \;\; f_\pi (X_\pi + Y_\pi) \left[(m_{\Lambda_b} - m_p) f_1 + f_3 \frac{m_\pi^2}{m_{\Lambda_b}} \right] \, ,$$

$$b_\pi = f_\pi(X_\pi - Y_\pi)\left[(m_{\Lambda_b}+m_p)g_1 - g_3\frac{m_\pi^2}{m_{\Lambda_b}}\right] , \tag{26}$$

where

$$\begin{aligned} X_\pi &= \frac{G_F}{\sqrt{2}}\left[V_{ub}V_{ud}^* a_2 - \sum_{q=u,c,t} V_{qb}V_{qd}^*(a_4^q + a_{10}^q) - V_{tb}V_{td}^* a_d(1+\frac{2E_\pi}{m_b})\right] , \\ Y_\pi &= -\frac{G_F}{\sqrt{2}}\left[\sum_{q=u,c,t} V_{qb}V_{qd}^*(a_6^q + a_8^q) + \frac{5}{4}V_{tb}V_{td}^* a_d\right]\chi_\pi , \end{aligned} \tag{27}$$

with

$$\chi_\pi = \frac{2m_\pi^2}{(m_d+m_u)m_b} . \tag{28}$$

Finally, we consider the decay $\Lambda_b \to \Lambda\eta(\eta')$ [12], which is dominated by a colour-allowed $b \to s$ penguin transition (there is also a small colour-suppressed tree contribution). For the decay $\Lambda_b \to \Lambda\eta$ we get

$$\begin{aligned} a_\eta &= f_\pi(X_\eta + Y_\eta)\left[(m_{\Lambda_b}-m_\Lambda)f_1 + f_3\frac{m_\eta^2}{m_{\Lambda_b}}\right] , \\ b_\eta &= f_\pi(X_\eta - Y_\eta)\left[(m_{\Lambda_b}+m_\Lambda)g_1 - g_3\frac{m_\eta^2}{m_{\Lambda_b}}\right] , \end{aligned} \tag{29}$$

where

$$\begin{aligned} X_\eta &= \frac{G_F}{\sqrt{2}}\left[V_{ub}V_{us}^* a_1 r_1 - \sum_{q=u,c,t} V_{qb}V_{qs}^*(r_1 A_q + r_2 B_q)\right. \\ &\qquad \left. - V_{tb}V_{ts}^* r_2 a_d\left(1+\frac{2E_\eta}{m_b}\right)\right] , \\ Y_\eta &= -\frac{G_F}{\sqrt{2}}\left[\sum_{q=u,c,t} V_{qb}V_{qs}^*(a_6^q - \frac{1}{2}a_8^q) + \frac{5}{4}V_{tb}V_{ts}^* a_d\right] r_2\chi_\eta , \end{aligned} \tag{30}$$

with

$$\begin{aligned} A_q &= 2a_3^q - 2a_5^q - \frac{1}{2}a_7^q + \frac{1}{2}a_9^q , \\ B_q &= a_3^q + a_4^q - a_5^q + \frac{1}{2}a_7^q - \frac{1}{2}a_9^q - \frac{1}{2}a_{10}^q , \\ \chi_\eta &= \frac{m_\eta^2}{m_s m_b} . \end{aligned} \tag{31}$$

In the above, we have defined $r_1 = f_\eta^u/f_\pi$ and $r_2 = f_\eta^s/f_\pi$, with

$$\begin{aligned} i f_\eta^u p_\eta^\mu &= \langle\eta|\,\bar{u}\gamma^\mu(1-\gamma_5)u\,|0\rangle = \langle\eta|\,\bar{d}\gamma^\mu(1-\gamma_5)d\,|0\rangle \ , \\ i f_\eta^s p_\eta^\mu &= \langle\eta|\,\bar{s}\gamma^\mu(1-\gamma_5)s\,|0\rangle \ . \end{aligned} \tag{32}$$

The amplitude for $\Lambda_b \to \Lambda\eta'$ has the same form as Eq. (30) with the replacement $\eta \to \eta'$. Note that the polarization of the final-state Λ can be measured via its decay $\Lambda \to p\pi^-$.

The above analysis has been performed within the framework of factorization. Before turning to estimates of the size of the triple-product asymmetries, it is useful at this point to address the issue of nonfactorizable corrections. Nonfactorizable effects are known to be important for hyperon and charmed-baryon nonleptonic decays, but are expected to be negligible for non-leptonic Λ_b decays. An unambiguous signal for the presence of nonfactorizable effects would be the observation of the decay $\Lambda_b \to \Delta^+ K^-(\pi^-)$, $\Lambda_b \to \Sigma\eta(\eta')$, or $\Lambda_b \to \Sigma\phi$. This is because, for the factorizable contribution, the light diquark in the Λ_b baryon remains inert during the weak decay. Thus, since the light diquark is an isosinglet, and since strong interactions conserve isospin to a very good approximation, the above Λ_b decays are forbidden within factorization [13].

One way to estimate the size of nonfactorizable corrections is by using the pole model. In this model, one assumes that the nonfactorizable decay amplitude receives contributions primarily from one-particle intermediate states, and that these contributions then show up as simple poles in the decay amplitude. An example of intermediate single-particle states is the ground-state positive-parity baryons. Consider the decay $\Lambda_b \to pK^-$. One nonfactorizable contribution is described by the diagram in which there is a $\Lambda_b \to \Sigma^0$ weak transition through a W exchange, followed by the strong decay $\Sigma^0 \to pK^-$. The pole contribution to the parity-violating amplitude, a, in Eq. (3) is known to be small for charmed-baryon decays [14], and we assume this to be the case here as well. For the parity-conserving amplitude, b, in Eq. (3), we can then write

$$b_{nonfac} \sim V_{ub} V_{us}^* \frac{\langle \Sigma^0 | H_w | \Lambda_b \rangle}{m_{\Lambda_b} - m_{\Sigma_0}} g_{\Sigma^0 p K^-} \ , \tag{33}$$

where $g_{\Sigma^0 p K^-}$ is the strong-coupling vertex which will depend on the energy of the emitted kaon. We can use heavy-quark and flavour $SU(3)$ symmetry to set $\langle \Sigma^0 | H_w | \Lambda_b \rangle \sim \langle \Sigma^+ | H_w | \Lambda_c \rangle$. Writing the weak matrix element $\langle \Sigma^+ | H_w | \Lambda_c \rangle = \frac{G_F}{\sqrt{2}} m^3$, we obtain

$$\frac{b_{nonfac}}{b_{fac}} \sim \frac{m}{f_K} \frac{m^2}{(m_{\Lambda_b} - m_{\Sigma^0})(m_{\Lambda_b} + m_p)} g_{\Sigma^0 p K^-} \ , \tag{34}$$

where we have chosen the tree-level term for A_{fac}. Since the emitted kaon is hard and since the quarks inside it are energetic, the strong coupling $g_{\Sigma^0 p K^-} \sim \alpha_s(\mu \sim E_K \sim m_b)$. In other words, the offshell Σ^0 has to emit a hard gluon to create a $\bar{u}u$ pair to form the pK^- final state. The matrix element m can either be estimated using a model [14], or obtained from a fit to the charmed baryon decay $\Lambda_c \to \Sigma^0 \pi^+$ [15]. In both cases one obtains $m \sim 0.1 - 0.2$ GeV, so that, from Eq. (34), the nonfactorizable corrections are found to be tiny. Arguments for small nonfactorizable effects in Λ_b decays can also be made based on the total width calculations [16].

To summarize: for colour-allowed $\Lambda_b \to F_1 P$ decays, we find that the triple-product correlation $\text{Im}(ab^*) \vec{p}_{F_1} \cdot (\vec{s}_{F_1} \times \vec{s}_{\Lambda_b})$ can be nonzero. The next step is to calculate the size of the asymmetry $\mathcal{A}_T$ in Eq. (2) for the various decays.

We begin with $\Lambda_b \to pK^-$. For this decay, we use the expressions for a_K and b_K found in Eq. (22). We note that the f_3 (g_3) term is suppressed relative to the f_1 (g_1) term by a factor $m_K^2/m_{\Lambda_b}^2 \sim 0.01$, and so can be neglected (and similarly for the g_3 piece). Furthermore, we take $f_1 = g_1$ [Eq. (21)], in which case all dependence on this form factor cancels in A_T [Eq. (1)]. The quantities

a_K and b_K depend on the parameters of the Cabibbo-Kobayashi-Maskawa (CKM) matrix, whose values are taken to be

$$\rho = 0.17 \quad , \qquad \eta = 0.39 \ . \tag{35}$$

For the chiral enhancement term χ_K [Eq. (19)], we take $\chi_K = 1$.

In order to estimate the value of rescattering terms, one has to choose a value of q^2. We consider two possibilities:

$$\text{Model 1}: \ q^2 = \frac{m_b^2}{4} \quad , \qquad \text{Model 2}: \ q^2 = \frac{m_b^2}{2} \ . \tag{36}$$

Taking $m_c = 1.4$ GeV, $m_u = 6$ MeV and $m_b = 5$ GeV, and writing $c_4^c = |c_4^c| e^{i\delta^c}$, we find

$$\begin{aligned} \text{Model 1}: \ & |c_4^u| = 0.02 \ , \ |c_4^c| = 0.02 \ , \ \delta^c = 51^\circ \ , \\ \text{Model 2}: \ & |c_4^u| = 0.021 \ , \ |c_4^c| = 0.015 \ , \ \delta^c = 0 \ . \end{aligned} \tag{37}$$

(In accordance with CPT, we set the phase of c_4^u to zero [17].)

Before presenting the numerical analysis, it is useful to anticipate the results. Referring to Eqs. (6), (22) and (18), we expect the triple-product asymmetry to be of order

$$\frac{2\mathrm{Im}(a_K b_K^*)}{|a_K|^2 + |b_K|^2 + 2\mathrm{Re}(a_K b_K^*)} \simeq \frac{\mathrm{Im}(X_K Y_K^*)}{|X_K|^2} \simeq \frac{a_2 a_6 \eta \lambda^2}{a_2^2 \lambda^4 + a_4^2} = 24\% \ . \tag{38}$$

Of course, this is a back-of-the-envelope estimate, but it does indicate that we can expect a reasonably large asymmetry, even when the rescattering effects are included. The fundamental reason for this is the following: the triple product is due mainly to the interference of the $V_{ub}V_{us}^*$ piece of X_K (we refer to this as T, the "tree") and the $V_{tb}V_{ts}^*$ piece of Y_K (P, the "penguin"). Like any CP-violating quantity, the asymmetry will therefore be maximized when the two interfering amplitudes are of comparable size. A quick calculation of these two quantities in Model 2 above shows that $|T/P| = 0.35$. The two amplitudes are therefore similar in size, leading to the sizeable asymmetry estimate above.

We have performed the phase-space integration for $\Lambda_b \to pK^-$ using the computer program RAMBO. For Model 1, we find that $A_T = -20.8\%$ and $\bar{A}_T = +15.0\%$, leading to a T-violating asymmetry of $\mathcal{A}_T^{pK} = -17.9\%$. In Model 2, since the strong phase vanishes, one necessarily has $\bar{A}_T = -A_T$, and we find $\mathcal{A}_T^{pK} = -19.1\%$. (We note in passing that the rescattering effects are quite important. Without them, the asymmetry would be $\mathcal{A}_T^{pK} = -26.1\%$. Thus, their inclusion leads to a correction in the asymmetry of about 25%.) These numbers are all consistent with the estimate in Eq. (38). We therefore conclude that the SM predicts a sizeable triple-product asymmetry in the decay $\Lambda_b \to pK^-$. (Note that, since the estimate in Eq. (38) uses only the values of the Wilson coefficients and the CKM matrix elements, we expect a large asymmetry even if nonfactorizable contributions are present).

There is one digressionary remark which is worth making here. From the measurement of ϵ_K, the CP-violating parameter in the kaon system, we know that the product $B_K\eta$ is positive, where B_K is the kaon bag parameter and η is the CP-violating CKM parameter. It is usually assumed that $B_K > 0$, so that η is also positive, and the unitarity triangle points up. However, there is no experimental evidence yet that $B_K > 0$. The T-violating triple-product asymmetry in $\Lambda_b \to pK^-$ is proportional to $\eta\cos(\delta)$, where δ is a strong phase. If one assumes that $|\delta| < 90°$, which is strongly favoured theoretically, then the triple product asymmetry measures the sign of η. This provides a cross check to the information obtained from the kaon system.

Turning now to the decay $\Lambda_b \to p\pi^-$, we have applied this same analysis as above. Taking $m_d = m_u = 6$ MeV, we have $\chi_\pi = 0.65$. In this case, the tree amplitude T is larger than the penguin amplitude P, with $|P/T| = 0.08$.

Because these two interfering amplitudes are less comparable in size than was the case for $\Lambda_b \to pK^-$, we expect a correspondingly smaller asymmetry. This is indeed what is found. In Model 1, we have $A_T = 6.3\%$ and $\bar{A}_T = -4.5\%$, so that $\mathcal{A}_T^{p\pi} = 5.4\%$. Model 2 gives a similar asymmetry: $\mathcal{A}_T^{p\pi} = 5.6\%$.

For the decays $\Lambda_b \to \Lambda\eta$ and $\Lambda_b \to \Lambda\eta'$, we have to define the quark content and mixing of the physical η and η' mesons. We use the Isgur mixing [18]:

$$\langle\eta| = \frac{1}{\sqrt{2}}[N - S] \ , \qquad \langle\eta'| = \frac{1}{\sqrt{2}}[N + S] \ , \tag{39}$$

where $N = [\langle u\bar{u}| + \langle d\bar{d}|]/\sqrt{2}$ and $S = \langle s\bar{s}|$. $SU(3)$ symmetry then gives

$$f_\eta^u = f_\pi/2 \ , \quad f_\eta^s = -f_\pi/\sqrt{2} \ , \quad f_{\eta'}^u = f_\pi/2 \ , \quad f_{\eta'}^s = f_\pi/\sqrt{2} \ , \tag{40}$$

where $f_\pi = 131$ MeV. We also take $\chi_\eta = 0.6$ and $\chi_{\eta'} = 1.8$. For both decays the interfering amplitudes are very different in size: $|T/P| = 0.03$ and 0.01 for the $\Lambda\eta$ and $\Lambda\eta'$ final states, respectively. We can therefore expect to obtain tiny triple-product asymmetries, and this should hold even if nonfactorizable effects are present. For $\Lambda_b \to \Lambda\eta$ we have $\mathcal{A}_T^{\Lambda\eta} = 0.6\%$ (Model 1) or 0.9% (Model 2), while for $\Lambda_b \to \Lambda\eta'$, $\mathcal{A}_T^{\Lambda\eta'} = -0.6\%$ (Model 1) or -0.5% (Model 2). It is unlikely that such tiny asymmetries can be measured. However, this also suggests that these processes might be good areas to search for new physics [19].

We now turn to the decays $\Lambda_b \to F_1 V$. The general decay amplitude can be written as [20]

$$\begin{aligned} \mathcal{M}_V &= Amp(\Lambda_{F_1} \to BV) \\ &= \bar{u}_{F_1}\varepsilon_\mu^*\Big[(p_{\Lambda_b}^\mu + p_{F_1}^\mu)(a + b\gamma_5) + \gamma^\mu(x + y\gamma_5)\Big]u_{\Lambda_b} \ , \end{aligned} \tag{41}$$

where ε_μ^* is the polarization of the vector meson. In the rest frame of the Λ_b, we can write $p_V = (E_V, 0, 0, |\vec{p}|)$ and $p_{F_1} = (E_{F_1}, 0, 0, -|\vec{p}|)$. Thus, it is clear that $\varepsilon_V^* \cdot (p_{\Lambda_b} + p_{F_1})$ will be nonzero only for a longitudinally-polarized V. This will be important in what follows.

The calculation of $|\mathcal{M}_V|^2$ gives the following triple-product terms:

$$\begin{aligned} |\mathcal{M}_V|^2_{t.p.} &= 2\,\mathrm{Im}(ab^*)\,|\varepsilon_V \cdot (p_{\Lambda_b} + p_{F_1})|^2 \; \epsilon_{\mu\nu\rho\sigma} p^\mu_{F_1} s^\nu_{F_1} p^\rho_{\Lambda_b} s^\sigma_{\Lambda_b} \\ &\quad + 2\,\mathrm{Im}\,(xy^*)\,\epsilon_{\alpha\beta\mu\nu} \Big[\varepsilon_V \cdot s_{F_1} p^\alpha_{F_1} p^\beta_{\Lambda_b} s^\mu_{\Lambda_b} \varepsilon^\nu_V - \varepsilon_V \cdot p_{F_1} s^\alpha_{F_1} p^\beta_{\Lambda_b} s^\mu_{\Lambda_b} \varepsilon^\nu_V \\ &\qquad + \varepsilon_V \cdot s_{\Lambda_b} p^\alpha_{F_1} s^\beta_{F_1} \varepsilon^\mu_V p^\nu_{\Lambda_b} - \varepsilon_V \cdot p_{\Lambda_b} p^\alpha_{F_1} s^\beta_{F_1} \varepsilon^\mu_V s^\nu_{\Lambda_b}\Big] \\ &\quad + 2\,\varepsilon_V \cdot (p_{\Lambda_b} + p_{F_1})\,\epsilon_{\alpha\beta\mu\nu} \Big[\mathrm{Im}\,(ax^* + by^*)\,p^\alpha_{F_1} s^\beta_{F_1} p^\mu_{\Lambda_b} \varepsilon^\nu_V \\ &\qquad + m_{\Lambda_b} \mathrm{Im}\,(bx^* + ay^*)\,p^\alpha_{F_1} s^\beta_{F_1} s^\mu_{\Lambda_b} \varepsilon^\nu_V \\ &\qquad - \mathrm{Im}\,(ax^* - by^*)\,p^\alpha_{F_1} p^\beta_{\Lambda_b} s^\mu_{\Lambda_b} \varepsilon^\nu_V \\ &\qquad - m_{F_1} \mathrm{Im}\,(ay^* - bx^*)\,s^\alpha_{F_1} p^\beta_{\Lambda_b} s^\mu_{\Lambda_b} \varepsilon^\nu_V\Big] \; . \end{aligned} \tag{42}$$

Note that if we sum over the polarization of the vector meson, we essentially reproduce the results found for $\Lambda_b \to F_1 P$. That is, there is only one triple product, which takes the form $\epsilon_{\mu\nu\rho\sigma} p^\mu_{F_1} s^\nu_{F_1} p^\rho_{\Lambda_b} s^\sigma_{\Lambda_b}$.

As usual, we use factorization to calculate the coefficients a, b, x and y. Consider first the decay $\Lambda_b \to pK^{*-}$. We define the decay constant g_{K^*} as

$$m_{K^*} g_{K^*} \varepsilon^*_\mu = \langle K^* |\, \bar{s}\gamma_\mu u \,|0\rangle \; . \tag{43}$$

In general, factorization allows us to write

$$\begin{aligned} A(\Lambda_b \to pK^{*-}) &= m_{K^*} g_{K^*} \Big\{ \varepsilon^*_\mu \langle p|\, \bar{u}\gamma^\mu(1-\gamma_5) b \,|\Lambda_b\rangle\, X_{K^*} \\ &\quad + \varepsilon^*_\mu \langle p|\, \bar{u}\gamma^\mu(1+\gamma_5) b \,|\Lambda_b\rangle\, Y_{K^*} \\ &\quad + \varepsilon \cdot (p_{\Lambda_b} + p_p)\, q_\mu \,\langle p|\, \bar{u}\gamma^\mu(1-\gamma_5) b \,|\Lambda_b\rangle\, A_{K^*} \\ &\quad + \varepsilon \cdot (p_{\Lambda_b} + p_p)\, q_\mu \,\langle p|\, \bar{u}\gamma^\mu(1+\gamma_5) b \,|\Lambda_b\rangle\, B_{K^*} \Big\} \; . \end{aligned} \tag{44}$$

The coefficients X_{K^*}, Y_{K^*}, A_{K^*} and B_{K^*} can be calculated using the effective hamiltonian. As noted earlier, A_{K^*} and B_{K^*} are nonzero only for a longitudinally-polarized K^{*-}.

Consider first the operators O_1–O_{10}. Since all of these lead to K^{*-} matrix elements of the form in Eq. (43), none of them can contribute to A_{K^*} and B_{K^*}. Furthermore, one can show that none of these give $Y_{K^*} \neq 0$ either. For example, consider again the operator O_6, which led to $Y_K \neq 0$. Because $\langle K^{*-}|\, \bar{s}(1+\gamma_5)u\,|0\rangle = 0$, O_6 will not contribute to Y_{K^*}. Thus, within factorization, if we restrict ourselves only to the operators O_1–O_{10}, the only nonzero coefficient is X_{K^*}, which means that all triple products vanish, since there is only a single decay amplitude.

In order to generate triple products in $\Lambda_b \to pK^{*-}$, it is necessary to consider the dipole operator O_{11}, whose effective coefficient is rather small (Eq. (19): $a_d = 0.0021$). However, there is an important observation one can make. The contributions of O_{11} to Y_{K^*}, A_{K^*} and B_{K^*} all involve the tensor matrix element for K^{*-}, which we define as

$$-ig^T_{K^*}\left[\varepsilon^*_\mu p^{K^*}_\nu - \varepsilon^*_\nu p^{K^*}_\mu\right] = \langle K^*|\, \bar{s}\sigma_{\mu\nu}u\,|0\rangle \ . \qquad (45)$$

Now, in the rest frame of the Λ_b, we can write $p_{K^*} = (E_{K^*}, 0, 0, |\vec{p}_{K^*}|)$. In the heavy-quark limit, in which $E_{K^*} \gg m_{K^*}$, the longitudinal polarization vector can be written approximately as

$$\varepsilon^{\lambda=0}_\mu \simeq \frac{1}{m_{K^*}}\left(p^\mu_{K^*} + \frac{m^2_{K^*}}{2E_{K^*}}n^\mu\right), \qquad (46)$$

with $n^\mu = (-1, 0, 0, 1)$. From Eq. (45), we see that the piece of $\varepsilon^{\lambda=0}_\mu$ which is proportional to $p^\mu_{K^*}$ will not contribute to the matrix element. Thus, the contribution of O_{11} to A_{K^*} and B_{K^*} is small, $O(m_{K^*}/2E_{K^*})$. Combining this with the small value of a_d, one obtains $A_{K^*} \simeq B_{K^*} \approx 0$. Furthermore, the above expression implies that the value of Y_{K^*} for a longitudinally-polarized K^{*-} meson will be suppressed relative to that for a transversely-polarized K^{*-} by about $m_{K^*}/2E_{K^*} = 16\%$. As we will see, Y_{K^*} is already small for a transversely-polarized K^{*-}, so that $Y_{K^*} \approx 0$ for longitudinal polarization. Therefore, any triple products in $\Lambda_b \to pK^{*-}$ should be largest for a transversely-polarized K^{*-}, although we expect even these to be small.

Considering separately the longitudinal ($\lambda = 0$) and transverse ($\lambda =\perp$) polarizations of the final-state vector meson, we find

$$\begin{aligned}
X_{K^*}^{\lambda=\perp} &= \frac{G_F}{\sqrt{2}}\left[V_{ub}V_{us}^* a_2 - \sum_{q=u,c,t} V_{qb}V_{qs}^*(a_4^q + a_{10}^q) - V_{tb}V_{ts}^*\frac{a_d}{2}\right] , \\
X_{K^*}^{\lambda=0} &= \frac{G_F}{\sqrt{2}}\left[V_{ub}V_{us}^* a_2 - \sum_{q=u,c,t} V_{qb}V_{qs}^*(a_4^q + a_{10}^q) - V_{tb}V_{ts}^* a_d \left(1 + \frac{2E_{K^*}}{m_b}\right)\right] , \\
Y_{K^*}^{\lambda=\perp} &= \frac{G_F}{\sqrt{2}} V_{tb}V_{ts}^* a_d , \\
Y_{K^*}^{\lambda=0} &\approx 0 ,
\end{aligned} \tag{47}$$

where

$$z \equiv \frac{E_{K^*}\, g_{K^*}^T}{m_{K^*}\, g_{K^*}} . \tag{48}$$

For $g_{K^*} = 226$ MeV and $g_{K^*}^T = 160$ MeV [21], $z = 2.23$. To a good approximation, the quantities a, b, x and y of Eq. (41) can then be expressed as

$$\begin{aligned}
a_{K^*}^\lambda &= m_{K^*} g_{K^*} \frac{f_2}{m_{\Lambda_b}} [X_{K^*}^\lambda + zY_{K^*}^\lambda] , \\
b_{K^*}^\lambda &= -m_{K^*} g_{K^*} \frac{g_2}{m_{\Lambda_b}} [X_{K^*}^\lambda - zY_{K^*}^\lambda] , \\
x_{K^*}^\lambda &= m_{K^*} g_{K^*} [f_1 - \frac{m_p + m_{\Lambda_b}}{m_{\Lambda_b}} f_2][X_{K^*}^\lambda + zY_{K^*}^\lambda] , \\
y_{K^*}^\lambda &= -m_{K^*} g_{K^*} [g_1 + \frac{m_{\Lambda_b} - m_p}{m_{\Lambda_b}} g_2][X_{K^*}^\lambda - zY_{K^*}^\lambda] .
\end{aligned} \tag{49}$$

There are several points to be deduced from the above results. First, since $Y_{K^*}^{\lambda=0} \approx 0$, the coefficients a, b, x and y all have the same phase for a longitudinally-polarized K^*. Thus all triple products involving a longitudinal

K^* are expected to vanish. Furthermore, since $\varepsilon_V \cdot p_{\Lambda_b} = 0$ for a transversely-polarized K^*, most of the triple products in Eq. (42) are expected to vanish in the SM. The only potential nonzero triple-product correlations are

$$2\,\mathrm{Im}\,(xy^*)\,\epsilon_{\alpha\beta\mu\nu}\left[\varepsilon_{K^*}\cdot s_p\,p_p^\alpha p_{\Lambda_b}^\beta s_{\Lambda_b}^\mu \varepsilon_{K^*}^\nu + \varepsilon_{K^*}\cdot s_{\Lambda_b}\,p_p^\alpha s_p^\beta \varepsilon_{K^*}^\mu p_{\Lambda_b}^\nu\right] . \qquad (50)$$

Since these both require the measurement of all three spins, this result is consistent with the results of Ref. [3].

Second, and more importantly, both of these asymmetries only arise due to the interference between the (small) dipole term $Y_{K^*}^{\lambda=\perp}$ and the $V_{ub}V_{us}^*$ piece of $X_{K^*}^{\lambda=\perp}$. Thus, by analogy with Eq. (38), we estimate the size of the asymmetries to be roughly

$$\frac{2\mathrm{Im}(xy^*)}{|x|^2+|y|^2+2\mathrm{Re}(xy^*)} \simeq \frac{z\mathrm{Im}(X_{K^*}^{\lambda=\perp}Y_{K^*}^{\lambda=\perp *})}{|X_{K^*}^{\lambda=\perp}|^2} \simeq \frac{za_2a_d\eta\lambda^2}{a_2^2\lambda^4+a_4^2} \sim 2\% \ , \qquad (51)$$

which would be very difficult to measure. (Essentially, the asymmetry is reduced compared to that in $\Lambda_b \to pK^-$ by the factor $z|a_d/a_6| = 0.11$.) Furthermore, the decay $\Lambda_b \to pK^{*-}$ is dominated by the longitudinally-polarized K^{*-}; the rate for the production of a transversely-polarized K^{*-} is suppressed by the factor $(m_{K^*}/E_{K^*})^2 = 0.1$. Thus, even if the asymmetry were larger, it would still be difficult to detect, given the small rate.

We therefore conclude that *any* measurement of a sizeable triple-product asymmetry in the decay $\Lambda_b \to pK^{*-}$ is an unequivocal signal of new physics [19]. (As noted in the case of $\Lambda_b \to pK^-$ decay, if the measurement of the proton polarization is difficult, one can consider a final state with an excited nucleon such as $\Lambda_b \to N(1440)K^{*-}$. In this case, the polarization of the $N(1440)$ can be determined from its decay products. Alternatively, one can consider $\Xi_b \to \Sigma^+K^{*-}$, for which the above conclusions should also hold.)

The decay $\Lambda_b \to p\rho^-$ is similar to $\Lambda_b \to pK^{*-}$, and its amplitude can be obtained from Eqs. (41), (49) and (47) with the replacements $V_{is} \to V_{id}$ and

$K^* \to \rho$. However, here too the asymmetry is expected to be smaller than that in $\Lambda_b \to p\pi^-$ by the factor $z|a_d/a_6| = 0.11$, yielding an asymmetry of less than 1%. Finally, the pure penguin process $\Lambda_b \to \Lambda\phi$, is dominated by a single weak amplitude, so that all its triple-product asymmetries vanish [22]. (As mentioned earlier, the observation of the decay $\Lambda_b \to \Sigma\phi$ would indicate the existence of nonfactorizable contributions and the possible presence of a significant $V_{ub}V_{us}^*$ piece in the amplitude.)

To summarize, we have examined the predictions of the standard model for T-violating triple-product asymmetries in $\Lambda_b \to F_1F_2$ decays, where F_1 is a light spin-$\frac{1}{2}$ baryon, and F_2 is a pseudoscalar (P) or vector (V) meson. In $\Lambda_b \to F_1P$ decays, there is only a single triple product possible. In the rest frame of the Λ_b, it takes the form $\vec{p}_{F_1}\cdot(\vec{s}_{F_1}\times\vec{s}_{\Lambda_b})$, where $\vec{p}_{F_1}$ is the 3-momentum of the F_1, and $\vec{s}_{F_1}$ and $\vec{s}_{\Lambda_b}$) are the spins of the F_1 and Λ_b, respectively. On the other hand, in $\Lambda_b \to F_1V$ decays, since all three particles have a non-zero spin, there are several possible triple products.

Using factorization, we find the following results. First, for $\Lambda_b \to F_1P$ decays, the SM predicts a large asymmetry ($\sim$ 18%) only for $\Lambda_b \to pK^-$. This is due to the presence of the chiral enhancement term $m_K^2/(m_s+m_u)$ in the amplitude, which compensates the $1/m_b$ suppression. The asymmetry in $\Lambda_b \to p\pi^-$ is smaller ($\sim$ 5%), and for the decays $\Lambda_b \to \Lambda\eta, \Lambda\eta'$, it is less than 1%. Second, for $\Lambda_b \to F_1V$ decays with a transversely-polarized V, the asymmetries are quite small: for $\Lambda_b \to pK^{*-}$ and $\Lambda_b \to p\rho^-$ they are $O(1\%)$ and $< 1\%$, respectively. The asymmetries involving a longitudinally-polarized V are expected to be roughly 15% smaller than those for a transversely-polarized V, so that they are effectively unmeasurable. There are no asymmetries in $\Lambda_b \to \Lambda\phi$ since this decay is dominated by a single weak amplitude. The fact that, within the SM, the triple-product asymmetries in many decays are tiny suggests that this is a good area to search for new physics.

Acknowledgements: We thank S. Pakvasa and R. MacKenzie for helpful discussions. This work was financially supported by NSERC of Canada.

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CHAPITRE III: TROISIÈME ARTICLE

(Publié dans: Physical Review D 66, 094004 (2002).)

New-Physics Effects on Triple-Product Correlations in Λ_b Decays

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Abstract

We adopt an effective-lagrangian approach to compute the new-physics contributions to T-violating triple-product correlations in charmless Λ_b decays. We use factorization and work to leading order in the heavy-quark expansion. We find that the standard-model (SM) predictions for such correlations can be significantly modified. For example, triple products which are expected to vanish in the SM can be enormous ($\sim 50\%$) in the presence of new physics. By measuring triple products in a variety of Λ_b decays, one can diagnose which new-physics operators are or are not present. Our general results can be applied to any specific model of new physics by simply calculating which operators appear in that model.

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1 Introduction

The origin of CP violation remains one of the important open questions in particle physics. Within the standard model (SM), CP violation is due to the presence of phases in the Cabibbo-Kobayashi-Maskawa (CKM) quark mixing matrix. The B-factories BaBar and Belle have been built to test this: if the SM explanation is correct, we expect to observe large CP-violating rate asymmetries in B decays [1]. To date, one of the CP phases of the unitarity triangle has been measured: $\sin 2\beta = 0.78 \pm 0.08$ [2], which is consistent with the SM.

Although the main focus has been on rate asymmetries, there is another type of CP-violating signal which could potentially reveal the presence of physics beyond the SM. Triple-product correlations of the form $\vec{v}_1 \cdot (\vec{v}_2 \times \vec{v}_3)$, where each v_i is a spin or momentum, are odd under time reversal (T). Therefore, by the CPT theorem, these are also signals of CP violation. A nonzero triple-product correlation is signalled by a nonzero value of the asymmetry

$$A_T \equiv \frac{\Gamma(\vec{v}_1 \cdot (\vec{v}_2 \times \vec{v}_3) > 0) - \Gamma(\vec{v}_1 \cdot (\vec{v}_2 \times \vec{v}_3) < 0)}{\Gamma(\vec{v}_1 \cdot (\vec{v}_2 \times \vec{v}_3) > 0) + \Gamma(\vec{v}_1 \cdot (\vec{v}_2 \times \vec{v}_3) < 0)} , \tag{1}$$

where Γ is the decay rate for the process in question. However, there is a well-known caveat: strong phases can produce a nonzero value of A_T, even if the weak phases are zero (i.e. CP violation is not really present). Thus, to be sure that one is truly probing T and CP violation, one must compare the value of A_T with that of $\bar{A}_T$, which is the T-odd asymmetry measured in the CP-conjugate decay process.

Triple-product correlations can be measured in $B \to V_1 V_2$ decays, where V_1 and V_2 are vector mesons [3]. In the rest frame of the B, the triple product takes the form $\vec{p} \cdot (\varepsilon_1 \times \varepsilon_2)$, where $\vec{p}$ is the momentum of one of the final-state particles, and ε_i is the polarization of the V_i. One can also consider triple-product correlations in Λ_b decays. Since many such triple products involve the spin of the Λ_b, this means that, in contrast to B decays, one is sensitive to the spin of the b-quark [4], as it is expected to provide the dominant contribution to the spin of the Λ_b.

In a recent paper [5], we used factorization to study the SM predictions for triple products in charmless two-body Λ_b decays. We considered decays which are generated by the quark-level transitions $b \to s\bar{q}q$ or $b \to d\bar{q}q$. These decays take the form $\Lambda_b \to F_1 F_2$, where F_1 is a light spin-$\frac{1}{2}$ baryon, such as p, Λ, etc., and F_2 is a pseudoscalar (P) or vector (V) meson. There was only one decay in which there was a large effect: the triple-product asymmetry for $\Lambda_b \to pK^-$ was found to be 18%. For all other decays, the asymmetries are found to be at most at the percent level.

The fact that all these triple-product asymmetries are expected to be small in the SM suggests that this is a good area to look for physics beyond the SM. In this paper, we examine the effect of new physics on triple products in charmless Λ_b decays. In order to study this, we adopt an effective-lagrangian approach: we write down all possible dimension-six new-physics four-fermi operators at the quark level. Then, using factorization, we compute their contributions to the various triple-product correlations in Λ_b decays.

There are several advantages to this approach. First, we are able to establish which triple products can be significantly affected by the presence of new physics. Second, we can also determine specifically which new-physics operators contribute to these triple products. Thus, by measuring a number of different triple-product correlations, we may be able to diagnose which operators are or are not present. Finally, these operators include all possible models of new physics. Therefore one can apply our results to a specific model by simply calculating which new-physics operators appear in that model. We will give examples of this procedure.

The paper is organized as follows. In Section 2, we introduce the new-physics operators used in our analysis. We also give two examples of specific models which generate some of these operators: supersymmetry with R-parity breaking, and Z- and Z'-mediated flavour-changing neutral currents. We compute the contributions of the new-physics operators to triple-product correlations in Λ_b decays in Section 3. Here we retain only the leading term in the heavy-quark expansion since it is very unlikely that new physics contributes to subleading processes without affecting the leading-order processes. In Section 4, we estimate the size of the various triple products in the presence of new physics. By comparing triple products in $\Lambda_b \to F_1 P$ and $\Lambda_b \to F_1 V$ decays, we examine the "diagnostic power" of this approach, i.e. the extent to which one can determine which new-physics operators are present. We also show how our results can be applied to the specific models of new physics discussed previously. We conclude in Section 5.

2 New Physics

We are interested in charmless Λ_b decays, which are governed by the quark-level processes $b \to s\bar{q}q$ or $b \to d\bar{q}q$. In what follows we will concentrate on the $b \to s$ transitions; it is straightforward to adapt our analysis to the $b \to d$ case.

Taking into account the two different colour structures, as well as all possible Lorentz structures, there are a total of 20 dimension-six new-physics operators which contribute to each of the $b \to s\bar{q}q$ transitions, $q = u, d, s$. These can be written as

$$\mathcal{H}_{NP}^{q} = \sum_{A,B=L,R} \frac{4G_F}{\sqrt{2}} \left\{ f_{q,1}^{AB}\, \bar{s}_\alpha \gamma_A b_\beta\, \bar{q}_\beta \gamma_B q_\alpha + f_{q,2}^{AB}\, \bar{s}\gamma_A b\, \bar{q}\gamma_B q \right.$$
$$+ g_{q,1}^{AB}\, \bar{s}_\alpha \gamma^\mu \gamma_A b_\beta\, \bar{q}_\beta \gamma_\mu \gamma_B q_\alpha + g_{q,2}^{AB}\, \bar{s}\gamma^\mu \gamma_A b\, \bar{q}\gamma_\mu \gamma_B q$$
$$\left. + h_{q,1}^{AB}\, \bar{s}_\alpha \sigma^{\mu\nu} \gamma_A b_\beta\, \bar{q}_\beta \sigma_{\mu\nu} \gamma_B q_\alpha + h_{q,2}^{AB}\, \bar{s}\sigma^{\mu\nu} \gamma_A b\, \bar{q}\sigma_{\mu\nu} \gamma_B q \right\}, \quad (2)$$

where we have defined $\gamma_{R(L)} = \frac{1}{2}(1 \pm \gamma_5)$. Note: although we have written the tensor operators in the same compact form as the other operators, it should be noted that those with $\gamma_A \neq \gamma_B$ are identically zero. Thus, one can effectively set $h_{q,i}^{LR} = h_{q,i}^{RL} = 0$.

All models of new physics which contribute to $b \to s\bar{q}q$ will generate operators found in the above effective hamiltonian. These can arise at tree level (e.g. supersymmetry with R-parity breaking, Z- and Z'-mediated flavour-changing neutral currents, models with flavour-changing neutral scalars, etc.) or at loop level (e.g. minimal supersymmetry, left-right symmetric models, four generations, etc.) [6]. In some cases one will obtain operators of the form $\bar{q}\mathcal{O}b\, \bar{s}\mathcal{O}'q$, but one can perform a Fierz transformation to put them into the form of Eq. (2). Note that, in general, models of new physics do not lead directly to tensor operators ($h_{q,i}^{AB}$), since typically only vector or scalar particles are involved. However, such tensor operators can arise when other operators are Fierz-transformed into the above form, so they must be included in our analysis (the scalar operator $\bar{q}b\bar{s}q$ is such an example).

Because the new-physics operators are of dimension six, by dimensional analysis we expect them to be suppressed by a factor Λ^2, where Λ is the scale of new physics. However, with the normalization in Eq. (2), the suppression factor is only M_W^2. We therefore expect the size of the coefficients $f_{q,i}^{AB}$, $g_{q,i}^{AB}$ and $h_{q,i}^{AB}$ to be naturally of $O(M_W^2/\Lambda^2) \sim 10^{-2}$ for a new-physics scale of about 1 TeV.

Even so, these new-physics effects can be quite significant. In the SM, one finds only operators of the form $\bar{s}\gamma^{\mu}\gamma_L b\, \bar{q}\gamma_{\mu}\gamma_{L,R}q$, with both colour assignments. These operators are typically multiplied by one of two factors: either (i) the CKM matrix elements $V_{tb}V_{ts}^*$ times a Wilson coefficient of $O(10^{-2})$, or (ii) $V_{ub}V_{us}^*$ times a Wilson coefficient of $O(1)$. In either case, new-physics operators with coefficients of $O(10^{-2})$ would actually *dominate* over the SM contributions. (This is, in part, what allows us to put constraints on specific models of new physics.) The bottom line is that the new operators of Eq. (2) can contribute substantially to charmless Λ_b decays.

As noted above, by construction the effective hamiltonian of Eq. (2) includes all possible models of new physics. Of course, in a particular new-physics model, only a subset of the new operators will appear. Our general analysis can then be applied to that specific model by retaining only the coefficients of the nonzero operators. In order to show explicitly how this works, below we give two examples of such specific models.

2.1 Supersymmetry with R-parity breaking

In supersymmetric models, the R-parity of a field with spin S, baryon number B and lepton number L is defined to be

$$R = (-1)^{2S+3B+L} . \tag{3}$$

R is $+1$ for all the SM particles and -1 for all the supersymmetric particles. R-parity invariance is often imposed on the Lagrangian in order to maintain the separate conservation of baryon number and lepton number. Imposition of R-parity conservation has some important consequences: super particles must be produced in pairs in collider experiments and the lightest super particle (LSP) must be absolutely stable. The LSP therefore provides a good candidate for cold dark matter.

Despite the above-mentioned attractive features of R-parity conservation, this conservation is not dictated by any fundamental principle such as gauge invariance, so that there is no compelling theoretical motivation for it. The most general superpotential of the MSSM, consistent with $SU(3) \times SU(2) \times U(1)$ gauge symmetry and supersymmetry, can be written as

$$\mathcal{W} = \mathcal{W}_R + \mathcal{W}_{\not R} , \tag{4}$$

where $\mathcal{W}_R$ is the R-parity conserving piece, and $\mathcal{W}_{\not R}$ breaks R-parity. They are given by

$$\mathcal{W}_R = h_{ij} L_i H_2 E_j^c + h'_{ij} Q_i H_2 D_j^c + h''_{ij} Q_i H_1 U_j^c , \tag{5}$$

$$\mathcal{W}_{\not R} = \frac{1}{2}\lambda_{[ij]k} L_i L_j E_k^c + \lambda'_{ijk} L_i Q_j D_k^c + \frac{1}{2}\lambda''_{i[jk]} U_i^c D_j^c D_k^c + \mu_i L_i H_2 . \tag{6}$$

Here $L_i(Q_i)$ and $E_i(U_i, D_i)$ are the left-handed lepton (quark) doublet and lepton (quark) singlet chiral superfields, where i, j, k are generation indices and c denotes a charge conjugate field. $H_{1,2}$ are the chiral superfields representing the two Higgs doublets.

In the R-parity-violating superpotential [Eq. (6)], the λ and λ' couplings violate lepton number conservation, while the λ'' couplings violate baryon number conservation. $\lambda_{[ij]k}$ is antisymmetric in the first two indices and $\lambda''_{i[jk]}$ is antisymmetric in the last two indices. There are therefore 27 λ'-type couplings and 9 each of the λ and λ'' couplings. While it is theoretically possible to have both baryon-number and lepton-number violating terms in the Lagrangian, the non-observation of proton decay imposes very stringent conditions on their simultaneous presence [7]. One therefore assumes the existence of either L-violating couplings or B-violating couplings, but not both. The terms proportional to λ are not relevant to our present discussion and will not be considered further.

We begin with the B-violating couplings. The transition $b \to s\bar{u}u$ can be generated at tree level through the t-channel exchange of the d-squark, $\tilde{d}_R$, with strength proportional to $|\lambda''_{112}\lambda^{*''}_{113}|$. However, this product of couplings is already constrained to be $\sim 10^{-8}$ from $n - \bar{n}$ oscillations and double nucleon decay [8]. There are therefore no significant contributions to the new-physics operators of Eq. (2) corresponding to $q = u$.

Similarly, the antisymmetry of the B-violating couplings, $\lambda''_{i[jk]}$ in the last two indices implies that there are no operators that can generate the $b \to s\bar{s}s$ transition, so that all the operators in Eq. (2) vanish for $q = s$.

Finally, the operators that generate the $b \to s\bar{d}d$ transition are given by [9]

$$L_{eff} = \frac{\lambda''_{i12}\lambda''^{*}_{i13}}{4m^2_{\tilde{u}_i}}(\bar{d}_\alpha\gamma_\mu\gamma_R d_\alpha\, \bar{s}_\beta\gamma_\mu\gamma_R b_\beta - \bar{d}_\alpha\gamma_\mu\gamma_R d_\beta\, \bar{s}_\beta\gamma_\mu\gamma_R b_\alpha) \; . \tag{7}$$

Hence the only nonvanishing operators in Eq. (2) are

$$g^{RR}_{d,1} = -g^{RR}_{d,2} = -\frac{\sqrt{2}}{G_F}\frac{\lambda''_{i12}\lambda''^{*}_{i13}}{16m^2_{\tilde{u}_i}} \; . \tag{8}$$

As mentioned above, the constraint on $|\lambda''_{112}\lambda''^{*}_{113}|$ is at the 10^{-8} level. However, the constraint on $|\lambda''_{i12}\lambda''^{*}_{i13}|$, $i \neq 1$, comes only from the nonleptonic decay $B^- \to \bar{K}^0\pi^-$ [9], and is much weaker:

$$|\lambda''_{i12}\lambda''^{*}_{i13}|(i \neq 1) \leq 1.03 \times 10^{-2} \; , \tag{9}$$

where a squark mass $m_{\tilde{f}} = 100$ GeV has been assumed. We therefore find

$$|g^{RR}_{d,1}| = |g^{RR}_{d,2}| \leq 7.6 \times 10^{-3} \; . \tag{10}$$

We now turn to the L-violating couplings. In terms of four-component Dirac spinors, these are given by [10]

$$\begin{aligned} \mathcal{L}_{\lambda'} \;=\; -\lambda'_{ijk}\Big[&\tilde{\nu}^i_L\bar{d}^k_R d^j_L + \tilde{d}^j_L\bar{d}^k_R\nu^i_L + (\tilde{d}^k_R)^*(\bar{\nu}^i_L)^c d^j_L \\ &-\tilde{e}^i_L\bar{d}^k_R u^j_L - \tilde{u}^j_L\bar{d}^k_R e^i_L - (\tilde{d}^k_R)^*(\bar{e}^i_L)^c u^j_L\Big] + h.c. \end{aligned} \tag{11}$$

There are a variety of sources which bound the above couplings [8, 9]. For the sake of brevity we will only quote the bounds and not their sources. Assuming a commom sfermion mass of 100 GeV we find the most stringent bounds are

$$|\lambda'_{i12}\lambda'^{*}_{i13}|(i \neq 1) \leq 1.7 \times 10^{-3} \quad , \quad |\lambda'_{112}\lambda'^{*}_{113}| \leq 4.4 \times 10^{-4} \tag{12}$$

$$\begin{array}{ccc} |\lambda'_{111*}\lambda'^{*}_{132}| \leq 1.4 \times 10^{-4} \ , & |\lambda'^{*}_{211}\lambda'_{232}| \leq 4.7 \times 10^{-4} \ , & |\lambda'_{311*}\lambda'_{332}| \leq 4.7 \times 10^{-4} \\ |\lambda'_{111}\lambda'^{*}_{123}| \leq 2.2 \times 10^{-5} \ , & |\lambda'_{211}\lambda'^{*}_{223}| \leq 2.2 \times 10^{-3} \ , & |\lambda'_{311}\lambda'^{*}_{323}| \leq 2.2 \times 10^{-3} \\ |\lambda'_{131}\lambda'^{*}_{121}| \leq 8.2 \times 10^{-4} \ , & |\lambda'_{231}\lambda'^{*}_{221}| \leq 1.3 \times 10^{-3} \ , & |\lambda'_{331}\lambda'^{*}_{321}| \leq 1.3 \times 10^{-3} \\ |\lambda'_{132}\lambda'^{*}_{122}| \leq 1.2 \times 10^{-2} \ , & |\lambda'_{232}\lambda'^{*}_{222}| \leq 1.2 \times 10^{-1} \ , & |\lambda'_{332}\lambda'^{*}_{322}| \leq 2.3 \times 10^{-1} \\ |\lambda'_{122}\lambda'^{*}_{123}| \leq 1.8 \times 10^{-3} \ , & |\lambda'_{223}\lambda'^{*}_{222}| \leq 4.4 \times 10^{-2} \ , & |\lambda'_{322}\lambda'^{*}_{323}| \leq 2.7 \times 10^{-1} \ . \end{array}$$

There is a single contribution to the $b \to s\bar{u}u$ transition:

$$L_{eff} = -\frac{\lambda'_{i12}\lambda'^{*}_{i13}}{2m^2_{\tilde{e}_i}} \bar{u}_\alpha \gamma_\mu \gamma_L u_\beta \, \bar{s}_\beta \gamma_\mu \gamma_R b_\alpha \ . \tag{13}$$

Hence the only nonvanishing operator for $q = u$ in Eq. (2) is

$$g^{RL}_{u,1} = -\frac{\sqrt{2}}{G_F} \frac{\lambda'_{i12}\lambda'^{*}_{i13}}{8m^2_{\tilde{e}_i}} \ . \tag{14}$$

Using the bounds of Eq. (13), we find

$$|g^{RL}_{u,1}| \leq 2.6 \times 10^{-3} \ . \tag{15}$$

Turning now to the $b \to s\bar{d}d$ transition the relevant Lagrangian is

$$\begin{aligned} L_{eff} \ = \ & \frac{\lambda'_{i11}\lambda'^{*}_{i23}}{m^2_{\tilde{\nu}_i}} \bar{d}\gamma_L d_\beta \, \bar{s}\gamma_R b + \frac{\lambda'_{i32}\lambda'^{*}_{i11}}{m^2_{\tilde{\nu}_i}} \bar{d}\gamma_R d \, \bar{s}\gamma_L b \\ & - \frac{\lambda'_{i12}\lambda'^{*}_{i13}}{2m^2_{\tilde{\nu}_i}} \bar{d}_\alpha \gamma_\mu \gamma_L d_\beta \, \bar{s}_\beta \gamma_\mu \gamma_R b_\alpha - \frac{\lambda'_{i31}\lambda'^{*}_{i21}}{2m^2_{\tilde{\nu}_i}} \bar{d}_\alpha \gamma_\mu \gamma_R d_\beta \, \bar{s}_\beta \gamma_\mu \gamma_L b_\alpha \ . \end{aligned} \tag{16}$$

The nonvanishing operators in Eq. 2 are then

$$f_{d,2}^{LR} = \frac{\sqrt{2}}{G_F}\frac{\lambda'_{i32}\lambda'^*_{i11}}{4m^2_{\tilde{\nu}_i}} \quad , \quad f_{d,2}^{RL} = \frac{\sqrt{2}}{G_F}\frac{\lambda'_{i11}\lambda'^*_{i23}}{4m^2_{\tilde{\nu}_i}} \ ,$$

$$g_{d,1}^{RL} = g_{u,1}^{RL*} \quad , \quad g_{d,1}^{LR} = -\frac{\sqrt{2}}{G_F}\frac{\lambda'_{i31}\lambda'^*_{i21}}{8m^2_{\tilde{\nu}_i}} \ , \tag{17}$$

with

$$|f_{d,2}^{LR}| \le 1.4 \times 10^{-3} \quad , \quad |f_{d,2}^{RL}| \le 6.6 \times 10^{-3} \ ,$$
$$|g_{d,1}^{RL}| \le 2.6 \times 10^{-3} \quad , \quad |g_{d,1}^{LR}| \le 2.0 \times 10^{-3} \ . \tag{18}$$

Finally, turning to the $b \to s\bar{s}s$ transition, the relevant Lagrangian is

$$L_{eff} = \frac{\lambda'_{i32}\lambda'^*_{i22}}{m^2_{\tilde{\nu}_i}}\bar{s}\gamma_R s\, \bar{s}\gamma_L b + \frac{\lambda'_{i22}\lambda'^*_{i23}}{m^2_{\tilde{\nu}_i}}\bar{s}\gamma_L s\, \bar{s}\gamma_R b \ , \tag{19}$$

allowing the identification

$$f_{s,2}^{LR} = \frac{\sqrt{2}}{G_F}\frac{\lambda'_{i32}\lambda'^*_{i22}}{4m^2_{\tilde{\nu}_i}} \quad , \quad f_{s,2}^{RL} = \frac{\sqrt{2}}{G_F}\frac{\lambda'_{i22}\lambda'^*_{i23}}{4m^2_{\tilde{\nu}_i}} \ , \tag{20}$$

with

$$|f_{s,2}^{LR}| \le 0.7 \quad , \quad |f_{s,2}^{RL}| \le 0.8 \ . \tag{21}$$

2.2 Z- and Z'-mediated FCNC's

In these models, one introduces an additional vector-singlet charge $-1/3$ quark h, as is found in E_6 grand unified theories, and allows it to mix with the ordinary down-type quarks d, s and b. Since the weak isospin of the exotic quark is different from that of the ordinary quarks, flavor-changing neutral currents (FCNC's) involving the Z are induced [11]. The $Zb\bar{s}$ FCNC coupling, which is of interest to us here, is parametrized by the independent parameter U^Z_{sb}:

$$\mathcal{L}^Z_{FCNC} = -\frac{g}{2\cos\theta_W} U^Z_{sb}\, \bar{s}_L \gamma^\mu b_L Z_\mu \ . \tag{22}$$

Note that it is only the mixing between the left-handed components of the ordinary and exotic quarks which is responsible for the FCNC: since s_R, b_R and h_R all have the same $SU(2)_L \times U(1)_Y$ quantum numbers, their mixing cannot generate flavour-changing couplings of the Z. Models with Z-mediated FCNC's will therefore generate the $g^{LL}_{q,2}$ and $g^{LR}_{q,2}$ new-physics operators Eq. (2). These are the same operators that appear in the SM, so that this model does not generate new operators. (That is, these are effectively new contributions to the electroweak penguin operators of the SM.)

The strongest constraint on U^Z_{sb} comes from the measurement of $B(B \to \ell^+\ell^- X)$. The most recent result from BELLE gives [12]

$$B(B \to X_s e^+ e^-) \leq 1.01 \times 10^{-5} \ , \tag{23}$$

leading to the constraint

$$|U^Z_{sb}| \leq 7.6 \times 10^{-4} \ . \tag{24}$$

With this constraint, it is straightforward to compute the maximal size of the couplings $g^{LL}_{q,2}$ and $g^{LR}_{q,2}$. We find

$$\begin{aligned} |g^{LL}_{u,2}| \leq 2.7 \times 10^{-4} \quad &, \quad |g^{LR}_{u,2}| \leq 1.1 \times 10^{-4} \ , \\ |g^{LL}_{d,2}| = |g^{LL}_{s,2}| \leq 3.2 \times 10^{-4} \quad &, \quad |g^{LR}_{d,2}| = |g^{LR}_{s,2}| \leq 6.1 \times 10^{-5} \ . \end{aligned} \tag{25}$$

These couplings are therefore comparable in size to those of the SM.

Of course, since no new operators are generated in this scenario, and since the new-physics effects are about the same size as in the SM, one does not expect large deviations from the SM predictions due to Z-mediated FCNC's.

However, models of new physics which contain exotic fermions also predict, in general, the existence of additional neutral Z' gauge bosons. If the s-, b- and h-quarks have different quantum numbers under the new $U(1)$ symmetry, their mixing will induce FCNC's due to Z' exchange [13].

In general, as was the case for Z-mediated FCNC's, such flavour-changing couplings will be constrained by the measurement of $B(B \to \ell^+\ell^- X)$. However, if the Z' is leptophobic, i.e. it does not couple to charged leptons, one can evade the constraints due to Eq. (23). Such models were considered in Ref. [14]. In this case, it is the mixing of the right-handed components of the ordinary and exotic quarks which is most important, and we parametrize the flavour-changing $Z'b\bar{s}$ coupling as

$$\mathcal{L}^{Z'}_{FCNC} = -\frac{g}{2\cos\theta_W} U^{Z'}_{sb}\, \bar{s}_R \gamma^\mu b_R Z'_\mu \ . \tag{26}$$

Thus, these models will generate new operators. In particular, the coefficients $g^{RL}_{q,2}$ and $g^{RR}_{q,2}$ will be nonzero.

Even though the Z' is leptophobic, there are constraints on $U^{Z'}_{sb}$ coming from the ALEPH limit $B(b \to s\nu\bar{\nu}) \le 6.4 \times 10^{-4}$ [15]. In addition, in realistic models, leptophobia is realized only approximately – there will always be threshold effects which produce a small coupling of the Z' to charged leptons, in which case there are constraints from Eq. (23). The constraints from both of these sources turn out to be similar in size, and lead to [14]

$$\left|U^{Z'}_{sb}\right| \frac{M_Z^2}{M_{Z'}^2} \lesssim 6 \times 10^{-3} \ . \tag{27}$$

With this constraint, one can estimate how large the new-physics coefficients can be. One finds

$$|g^{AB}_{q,2}| \le (1-2) \times 10^{-3} \ , AB = RR, RL \ , q = u, d, s \ , \tag{28}$$

which is about an order of magnitude larger than the coefficients of Eq. (25). Thus, Z'-mediated FCNC's can contribute significantly to charmless hadronic

Λ_b decays, and can lead to significant deviations from the SM predictions for triple products in such processes.

3 Triple Products

In this section, we compute the contributions from the new-physics operators to triple-product correlations in Λ_b decays. In all cases, we retain only the leading term in the heavy-quark expansion, and neglect terms of order m/m_{Λ_b}, where m is the mass of the light meson. The main reason is that it is very unlikely that new physics will contribute at subleading order, but not at leading order. Indeed, as we will see, this situation can arise only in fine-tuned scenarios. A secondary reason is that the subleading terms are quite a bit smaller, e.g. $m_{K^*}/m_{\Lambda_b} \sim 15\%$.

We begin with a review of the results of the SM.

3.1 SM Results

In this subsection, we summarize the predictions of the SM for triple products in Λ_b decays. The discussion is somewhat cursory, and we refer to the reader to Ref. [5] for more details.

We first consider the decay $\Lambda_b \to F_1 P$, whose amplitude can be written generally as

$$\mathcal{M}_P = A(\Lambda_b \to F_1 P) = i\bar{u}_{F_1}(a + b\gamma_5)u_{\Lambda_b} \ . \qquad (29)$$

The calculation of $|\mathcal{M}_P|^2$ yields a single triple-product term:

$$\mathrm{Im}(ab^*)\epsilon_{\mu\nu\rho\sigma}p^{\mu}_{F_1}s^{\nu}_{F_1}p^{\rho}_{\Lambda_b}s^{\sigma}_{\Lambda_b} \ , \qquad (30)$$

where p_i^μ and s_i^μ are the 4-momentum and polarization of particle i. In the rest frame of the Λ_b, this takes the form $\vec{p}_{F_1} \cdot (\vec{s}_{F_1} \times \vec{s}_{\Lambda_b})$.

Within factorization, one can write

$$\begin{aligned} A(\Lambda_b \to F_1 P) &= \sum_{O,O'} \langle P|\, O\,|0\rangle \, \langle F_1|\, O'\,|\Lambda_b\rangle \\ &= i f_P q^\mu \, \langle F_1|\, \bar{q}_1 \gamma_\mu (1-\gamma_5) b\,|\Lambda_b\rangle \, X_P \\ &\qquad + i f_P q^\mu \, \langle F_1|\, \bar{q}_1 \gamma_\mu (1+\gamma_5) b\,|\Lambda_b\rangle \, Y_P \ , \end{aligned} \tag{31}$$

where we have defined the pseudoscalar decay constant f_P as

$$i f_P q^\mu = \langle P|\, \bar{q}_2 \gamma^\mu (1-\gamma_5) q_3\,|0\rangle \ , \tag{32}$$

where $q^\mu \equiv p^\mu_{\Lambda_b} - p^\mu_{F_1}$ is the four-momentum transfer. The key point here is that, in order to obtain a nonzero triple-product correlation, one must have two interfering amplitudes, i.e. X_P and Y_P must both be nonzero, and must have a relative weak phase. Furthermore, the triple product will be large only if X_P and Y_P are of similar size.

One can also show that the parameters a and b of Eq. (29) can be written as

$$\begin{aligned} a &= f_P (X_P + Y_P)(m_{\Lambda_b} - m_{F_1}) f_1 \ , \\ b &= f_P (X_P - Y_P)(m_{\Lambda_b} + m_{F_1}) g_1 \ , \end{aligned} \tag{33}$$

where we have dropped terms of $O(m_P/m_{\Lambda_b})$, and f_1 and g_1 are Lorentz-invariant form factors:

$$\begin{aligned} \langle F_1|\, \bar{q}_1 \gamma^\mu b\,|\Lambda_b\rangle &= \bar{u}_{F_1} \left[f_1 \gamma^\mu + i \frac{f_2}{m_{\Lambda_b}} \sigma^{\mu\nu} q_\nu + \frac{f_3}{m_{\Lambda_b}} q^\mu \right] u_{\Lambda_b} \\ \langle F_1|\, \bar{q}_1 \gamma^\mu \gamma_5 b\,|\Lambda_b\rangle &= \bar{u}_{F_1} \left[g_1 \gamma^\mu + i \frac{g_2}{m_{\Lambda_b}} \sigma^{\mu\nu} q_\nu + \frac{g_3}{m_{\Lambda_b}} q^\mu \right] \gamma_5 u_{\Lambda_b} \ . \end{aligned} \tag{34}$$

From Eqs. (30) and (33), we therefore see explicitly that the triple product in $\Lambda_b \to F_1 P$ is proportional to $\mathrm{Im}(ab^*) \sim \mathrm{Im}(X_P Y_P^*)$.

In the SM, there is only one class of decays which is expected to show a significant effect [5]: the triple-product correlation for $\Lambda_b \to pK^-$ is found to be $\sim 18\%$. For decays such as $\Lambda_b \to \Lambda\eta$, $\Lambda_b \to \Lambda\eta'$ and $\Lambda_b \to n\bar{K}^0$, the triple product is less than 1%. The fundamental reason for this is that $\Lambda_b \to pK^-$ is governed by the quark-level transition $b \to s\bar{u}u$, which has both a tree and a penguin contribution, whereas the other decays are dominated by the $b \to s$ penguin amplitude. Thus, for $\Lambda_b \to \Lambda\eta, \Lambda\eta', n\bar{K}^0$, there is essentially only a single decay amplitude, which precludes any CP- and T-violating effects.

In the above discussion, the size of the triple products has been estimated within factorization. However, it is well-known that nonfactorizable effects can be important in Λ_c decays. For example, the decay $\Lambda_c \to \Sigma^+\phi$ has been observed [16], and this can only proceed via a (nonfactorizable) W-exchange diagram. This then begs the question of whether nonfactorizable effects might be important in Λ_b decays. In fact, the answer is that Λ_b decays are *not* expected to be significantly affected by such effects. In Ref. [17], it was found that the W-exchange contributions to inclusive Λ_b decays are suppressed relative to those in Λ_c decays by $O(m_c/m_b)^3$. This implies that even for exclusive decays such nonfactorizable W-exchange terms are expected to be small. This was confirmed in Ref. [5]: the W-exchange contributions to $\Lambda_b \to pK^-$ were estimated using a pole model, and the ratio of nonfactorizable to factorizable contributions was found to be tiny. For these reasons, here and below we ignore all nonfactorizable effects in Λ_b decays.

Turning to $\Lambda_b \to F_1 V$, the general decay amplitude can be written as [18]

$$\begin{aligned} \mathcal{M}_V &= Amp(\Lambda_{F_1} \to BV) \\ &= \bar{u}_{F_1}\varepsilon^*_\mu \left[(p^\mu_{\Lambda_b} + p^\mu_{F_1})(a + b\gamma_5) + \gamma^\mu(x + y\gamma_5)\right] u_{\Lambda_b} , \qquad (35) \end{aligned}$$

where ε^*_μ is the polarization of the vector meson. In calculating $|\mathcal{M}_V|^2$, one finds many triple-product terms:

$$
\begin{aligned}
|\mathcal{M}_V|^2_{t.p.} = \; & 2\,\mathrm{Im}(ab^*)\,|\varepsilon_V\cdot(p_{\Lambda_b}+p_{F_1})|^2\,\epsilon_{\mu\nu\rho\sigma}p^{\mu}_{F_1}s^{\nu}_{F_1}p^{\rho}_{\Lambda_b}s^{\sigma}_{\Lambda_b} \\
& +2\,\mathrm{Im}\,(xy^*)\,\epsilon_{\alpha\beta\mu\nu}\Big[\varepsilon_V\cdot s_{F_1}p^{\alpha}_{F_1}p^{\beta}_{\Lambda_b}s^{\mu}_{\Lambda_b}\varepsilon^{\nu}_V-\varepsilon_V\cdot p_{F_1}s^{\alpha}_{F_1}p^{\beta}_{\Lambda_b}s^{\mu}_{\Lambda_b}\varepsilon^{\nu}_V \\
& \qquad +\varepsilon_V\cdot s_{\Lambda_b}p^{\alpha}_{F_1}s^{\beta}_{F_1}\varepsilon^{\mu}_V p^{\nu}_{\Lambda_b}-\varepsilon_V\cdot p_{\Lambda_b}p^{\alpha}_{F_1}s^{\beta}_{F_1}\varepsilon^{\mu}_V s^{\nu}_{\Lambda_b}\Big] \\
& +2\,\varepsilon_V\cdot(p_{\Lambda_b}+p_{F_1})\,\epsilon_{\alpha\beta\mu\nu}\Big[\mathrm{Im}\,(ax^*+by^*)\,p^{\alpha}_{F_1}s^{\beta}_{F_1}p^{\mu}_{\Lambda_b}\varepsilon^{\nu}_V \\
& \qquad +m_{\Lambda_b}\mathrm{Im}\,(bx^*+ay^*)\,p^{\alpha}_{F_1}s^{\beta}_{F_1}s^{\mu}_{\Lambda_b}\varepsilon^{\nu}_V \\
& \qquad -\mathrm{Im}\,(ax^*-by^*)\,p^{\alpha}_{F_1}p^{\beta}_{\Lambda_b}s^{\mu}_{\Lambda_b}\varepsilon^{\nu}_V \\
& \qquad -m_{F_1}\mathrm{Im}\,(ay^*-bx^*)\,s^{\alpha}_{F_1}p^{\beta}_{\Lambda_b}s^{\mu}_{\Lambda_b}\varepsilon^{\nu}_V\Big]\;. \qquad (36)
\end{aligned}
$$

Similar to $\Lambda_b \to F_1 P$ decays, using factorization, one can write

$$
\begin{aligned}
A(\Lambda_b\to F_1V) = \; & m_V g_V\Big\{\varepsilon^*_\mu\,\langle F_1|\,\bar{q}_1\gamma^\mu(1-\gamma_5)b\,|\Lambda_b\rangle\,X_V \\
& +\varepsilon^*_\mu\,\langle F_1|\,\bar{q}_1\gamma^\mu(1+\gamma_5)b\,|\Lambda_b\rangle\,Y_V \qquad (37) \\
& +\varepsilon\cdot(p_{\Lambda_b}+p_{F_1})\,q_\mu\,\langle F_1|\,\bar{q}_1\gamma^\mu(1-\gamma_5)b\,|\Lambda_b\rangle\,\frac{A_V}{m^2_{\Lambda_b}} \\
& +\varepsilon\cdot(p_{\Lambda_b}+p_{F_1})\,q_\mu\,\langle F_1|\,\bar{q}_1\gamma^\mu(1+\gamma_5)b\,|\Lambda_b\rangle\,\frac{B_V}{m^2_{\Lambda_b}}\Big\}\;,
\end{aligned}
$$

where the decay constant g_V has been defined as

$$
m_V g_V \varepsilon^*_\mu = \langle V|\,\bar{q}_2\gamma_\mu q_3\,|0\rangle\;. \qquad (38)
$$

(Above we have included explicit factors of $m^2_{\Lambda_b}$ so that A_V and B_V have the same dimensions as X_V and Y_V. This differs from Ref. [5]. Also, note that, since the magnitudes of p_{Λ_b}, p_{F_1} and q are all of order m_{Λ_b}, the A_V and B_V operators are not apriori smaller than the X_V and Y_V operators.) Hence, using factorization, the quantities a, b, x and y of Eq. (36) can be expressed as

$$
\begin{aligned}
a_V^\lambda &= m_V g_V \left[\frac{f_2}{m_{\Lambda_b}}(X_V^\lambda + Y_V^\lambda) + f_1 \frac{m_{\Lambda_b} - m_{F_1}}{m_{\Lambda_b}^2}(A_V^\lambda + B_V^\lambda)\right] , \\
b_V^\lambda &= m_V g_V \left[-\frac{g_2}{m_{\Lambda_b}}(X_V^\lambda - Y_V^\lambda) + g_1 \frac{m_{\Lambda_b} + m_{F_1}}{m_{\Lambda_b}^2}(A_V^\lambda - B_V^\lambda)\right] , \\
x_V^\lambda &= m_V g_V [f_1 - \frac{m_{\Lambda_b} + m_{F_1}}{m_{\Lambda_b}} f_2][X_V^\lambda + Y_V^\lambda] , \\
y_V^\lambda &= -m_V g_V [g_1 + \frac{m_{\Lambda_b} - m_{F_1}}{m_{\Lambda_b}} g_2][X_V^\lambda - Y_V^\lambda] , \qquad (39)
\end{aligned}
$$

where λ denotes the polarization of the final-state V, and we have again dropped subleading terms of $O(m_V/m_{\Lambda_b})$. If any two of the four terms in Eq. (37) above have a relative weak phase, their interference can lead to triple products.

In the SM, one finds that $A_V \simeq B_V \approx 0$, and that $Y_V \approx 0$ for a longitudinally-polarized V. For a transversely-polarized V, Y_V can be nonzero, but is still quite small. Thus, $\Lambda_b \to F_1 V$ decays are dominated by a single amplitude (the X_V term in Eq. (37) above), so that triple products in such decays are expected to be tiny. Specifically, one finds [5] that the triple-product asymmetry in $\Lambda_b \to pK^{*-}$ is $O(1\%)$ for a transversely-polarized K^{*-}, while for a longitudinally-polarized K^{*-} the asymmetry is $\ll 1\%$. For $\Lambda_b \to \Lambda\phi$ [19] and $\Lambda_b \to n\bar{K}^{*0}$, the asymmetries essentially vanish since these decays are dominated by a single weak decay amplitude (the $b \to s$ penguin).

3.2 $\Lambda_b \to F_1 P$: New Physics

We begin by considering the new-physics contributions to triple-product correlations in $\Lambda_b \to pK^-$ decays. Although this process is governed by the quark transition $b \to s\bar{u}u$, one still has to perform Fierz transformations on the operators in Eq. (2) to put them in a form appropriate for this decay. Using the relations

$$\begin{aligned}
if_K q^\mu &= \langle K|\,\bar{s}\gamma^\mu(1-\gamma_5)u\,|0\rangle \ , \\
\langle K|\,\bar{s}(1\pm\gamma_5)u\,|0\rangle &= \mp\frac{if_K m_K^2}{m_s+m_u} \ , \\
\langle p|\,\bar{u}(1\pm\gamma_5)b\,|\Lambda_b\rangle &= \frac{q^\mu}{m_b}\langle p|\,\bar{u}\gamma_\mu(1\mp\gamma_5)b\,|\Lambda_b\rangle \ ,
\end{aligned} \tag{40}$$

we find that the new-physics contributions to X_K and Y_K [Eq. (31)] are

$$\begin{aligned}
X_K^{NP} &= \frac{G_F}{\sqrt{2}}\left[\frac{1}{4}a_{u,1}^{RR}\chi_K + \frac{1}{2}a_{u,1}^{LR} + b_{u,1}^{LL} - b_{u,1}^{RL}\chi_K + 3c_{u,1}^{RR}\chi_K\right] , \\
Y_K^{NP} &= \frac{G_F}{\sqrt{2}}\left[-\frac{1}{4}a_{u,1}^{LL}\chi_K - \frac{1}{2}a_{u,1}^{RL} - b_{u,1}^{RR} + b_{u,1}^{LR}\chi_K - 3c_{u,1}^{LL}\chi_K\right] .
\end{aligned} \tag{41}$$

where

$$\chi_K \equiv \frac{2m_K^2}{(m_s+m_u)m_b} \ , \tag{42}$$

and we have defined

$$a_{q,1}^{AB} \equiv f_{q,1}^{AB} + \frac{1}{N_c}f_{q,2}^{AB} \ , \quad b_{q,1}^{AB} \equiv g_{q,1}^{AB} + \frac{1}{N_c}g_{q,2}^{AB} \ , \quad c_{q,1}^{AB} \equiv h_{q,1}^{AB} + \frac{1}{N_c}h_{q,2}^{AB} \ . \tag{43}$$

Note that we can obtain Y_K^{NP} from X_K^{NP}, up to an overall minus sign, simply by changing the chiralities $L \leftrightarrow R$.

As discussed in the previous subsection, within the SM the triple-product correlation for $\Lambda_b \to pK^-$ is expected to be large, $\sim 18\%$. However, since the new-physics operators of Eq. (2) can contribute to both X_K and Y_K, this prediction can easily be modified.

We now turn to the decay $\Lambda_b \to \Lambda\eta(\eta')$, which receives contributions from all three quark-level processes $b \to s\bar{q}q$, $q = u, d, s$. The calculation is similar to that above. For $\Lambda_b \to \Lambda\eta$ we find

$$
\begin{aligned}
X_\eta^{NP} &= \frac{G_F}{\sqrt{2}}\left[x_u + x_d + x_s + 3\,c_{s,1}^{RR}\chi_{\eta_s}\right], \\
x_{u(d)} &= r_{u(d)}\left[\frac{1}{2}(a_{u(d),2}^{RL} - a_{u(d),2}^{RR})\chi_{\eta_{u(d)}} + (b_{u(d),2}^{LL} - b_{u(d),2}^{LR})\right], \\
x_s &= r_s\left[\frac{1}{2}(a_{s,2}^{RL} - a_{s,2}^{RR} + \frac{1}{2}a_{s,1}^{RR} - 2b_{s,1}^{RL})\chi_{\eta_s}\right. \\
&\qquad \left. + (b_{s,1}^{LL} + b_{s,2}^{LL} - b_{s,2}^{LR} + \frac{1}{2}a_{s,1}^{LR})\right], \qquad (44) \\
Y_\eta^{NP} &= \frac{G_F}{\sqrt{2}}\left[y_u + y_d + y_s - 3\,c_{s,1}^{LL}\chi_{\eta_s}]\right], \\
y_{u(d)} &= r_{u(d)}\left[\frac{1}{2}(-a_{u(d),2}^{LR} + a_{u(d),2}^{LL})\chi_{\eta_{u(d)}} + (-b_{u(d),2}^{RR} + b_{u(d),2}^{RL})\right], \\
y_s &= r_s\left[\frac{1}{2}(-a_{s,2}^{LR} + a_{s,2}^{LL} - \frac{1}{2}a_{s,1}^{LL} + 2b_{s,1}^{LR})\chi_{\eta_s}\right. \\
&\qquad \left. + (-b_{s,1}^{RR} - b_{s,2}^{RR} + b_{s,2}^{RL} - \frac{1}{2}a_{s,1}^{RL})\right],
\end{aligned}
$$

with

$$\chi_{\eta_{u,d,s}} = \frac{m_\eta^2}{m_{u,d,s}m_b}\,, \qquad (45)$$

and

$$a_{q,2}^{AB} \equiv f_{q,2}^{AB} + \frac{1}{N_c}f_{q,1}^{AB} \quad , \quad b_{q,2}^{AB} \equiv g_{q,2}^{AB} + \frac{1}{N_c}g_{q,1}^{AB}\,. \qquad (46)$$

In the above, we have defined $r_{u,d,s} \equiv f_\eta^{u,d,s}/f_\pi$, where

$$
\begin{aligned}
if_\eta^u p_\eta^\mu &= \langle\eta|\,\bar{u}\gamma^\mu(1-\gamma_5)u\,|0\rangle = \langle\eta|\,\bar{d}\gamma^\mu(1-\gamma_5)d\,|0\rangle\,, \\
if_\eta^s p_\eta^\mu &= \langle\eta|\,\bar{s}\gamma^\mu(1-\gamma_5)s\,|0\rangle\,. \qquad (47)
\end{aligned}
$$

The amplitude for $\Lambda_b \to \Lambda\eta'$ has the same form as Eq. (45) with the replacement $\eta \to \eta'$.

Finally, we consider the decay $\Lambda_b \to n\bar{K}^0$, which is related by isospin to $\Lambda_b \to pK^-$. This is a pure penguin decay, with $b \to s\bar{d}d$. This decay will be much difficult to detect experimentally. Nevertheless, we include it here for completeness. We find

$$\begin{aligned} X_{\bar{K}}^{NP} &= \frac{G_F}{\sqrt{2}}\left[\frac{1}{4}a_{d,1}^{RR}\chi_{\bar{K}} + \frac{1}{2}a_{d,1}^{LR} + b_{d,1}^{LL} - b_{d,1}^{RL}\chi_{\bar{K}} + 3c_{d,1}^{RR}\chi_{\bar{K}}\right], \\ Y_{\bar{K}}^{NP} &= \frac{G_F}{\sqrt{2}}\left[-\frac{1}{4}a_{d,1}^{LL}\chi_{\bar{K}} - \frac{1}{2}a_{d,1}^{RL} - b_{d,1}^{RR} + b_{d,1}^{LR}\chi_{\bar{K}} - 3c_{d,1}^{LL}\chi_{\bar{K}}\right], \end{aligned} \tag{48}$$

where

$$\chi_{\bar{K}} \equiv \frac{2m_K^2}{(m_s + m_d)m_b} . \tag{49}$$

For each of the decays $\Lambda_b \to \Lambda\eta$, $\Lambda_b \to \Lambda\eta'$ and $\Lambda_b \to n\bar{K}^0$, the triple product is tiny in the SM. This is due essentially to the fact that these decays are dominated by a single weak decay amplitude (the $b \to s$ penguin). However, this is no longer true in the presence of new physics; on the contrary, there may be several decay amplitudes. The new-physics operators may therefore lead to sizeable triple products in these decays.

3.3 $\Lambda_b \to F_1 V$: New Physics

We now examine the new-physics contributions to triple products in $\Lambda_b \to F_1 V$ decays. Before turning to specific decays, one can make some very general observations.

First, the amplitude for the production of a transversely-polarized vector boson V is suppressed relative to that for a longitudinally-polarized V by a factor m_V/E_V. Since $E_V \sim m_{\Lambda_b}/2$, this means that this production amplitude is subleading in the heavy-quark expansion, and can be neglected. In other words, in our analysis, we will assume the vector meson in the decay $\Lambda_b \to F_1 V$ to be essentially longitudinally polarized. As explained earlier, this is justified by the fact that it is very unlikely that the new physics will affect the production of a transversely-polarized V without also affecting that of a longitudinally-polarized V.

Second, in the rest frame of the Λ_b, we can write the 4-momentum of the final state vector meson as $q_\mu = (E_V, 0, 0, |\vec{p}_V|)$, so that the longitudinal polarization vector takes the form $\varepsilon_\mu^{\lambda=0} = (1/m_V)(|\vec{p}_V|, 0, 0, E_V)$. In the heavy-quark limit, $E_V \gg m_V$. Thus, in this limit, the longitudinal polarization vector can be written approximately as

$$\varepsilon_\mu^{\lambda=0} \simeq \frac{1}{m_V}\left(q_\mu + \frac{m_V^2}{2E_V} n_\mu\right), \tag{50}$$

with $n_\mu = (-1, 0, 0, 1)$. In other words, to leading order in the heavy-quark expansion, $\varepsilon_\mu^{\lambda=0}$ is proportional to q_μ. This has two important consequences.

Consider first the A_V amplitude of Eq. (37), which is one of the four amplitudes describing $\Lambda_b \to F_1 V$ decays:

$$m_V g_V \varepsilon \cdot (p_{\Lambda_b} + p_{F_1})\, q_\mu \, \langle F_1 |\, \bar{q}_1 \gamma^\mu (1-\gamma_5) b \,| \Lambda_b \rangle \, \frac{A_V}{m_{\Lambda_b}^2} \,. \tag{51}$$

Since $p_{F_1}^\mu = (E_{F_1}, 0, 0, -|\vec{p}|)$, one sees that $\varepsilon_V^* \cdot (p_{\Lambda_b} + p_{F_1})$ will be nonzero only for a longitudinally-polarized V. Now, writing the quark content of the V as $\bar{q}_2 q_3$, the operators which correspond to the V take the form $\bar{q}_2(1 \pm \gamma_5) q_3$, $\bar{q}_2 \gamma^\mu (1 \pm \gamma_5) q_3$ or $\bar{q}_2 \sigma^{\mu\nu}(1 \pm \gamma_5) q_3$. In calculating the V matrix elements, these yield

$$\begin{aligned}
\langle V |\, \bar{q}_2 (1 \pm \gamma_5) q_3 \,| 0 \rangle &= 0 \,, \\
\langle V |\, \bar{q}_2 \gamma_\mu (1 \pm \gamma_5) q_3 \,| 0 \rangle &= m_V g_V \varepsilon_\mu^* \,, \\
\langle V |\, \bar{q}_2 \sigma_{\mu\nu} q_3 \,| 0 \rangle &= -i g_V^T \left[\varepsilon_\mu^* q_\nu - \varepsilon_\nu^* q_\mu \right] .
\end{aligned} \tag{52}$$

Thus, we see that it is only the tensor matrix element which could potentially contribute to A_V. However, to leading order in the heavy-quark expansion, $\varepsilon_\mu^{\lambda=0} \sim q_\mu$, so that the tensor matrix element vanishes. Thus, we have $A_V =$

$O(m_V/m_{\Lambda_b})$, even in the presence of new-physics operators, and we neglect it. This argument applies also to the B_V amplitude of Eq. (37). (By comparison, X_V and Y_V are expected to be $O(1)$, i.e. leading order in the heavy-quark expansion.)

Neglecting the A_V and B_V terms, Eq. (39) reduces to

$$\begin{aligned} a_V^\lambda &= m_V g_V \frac{f_2}{m_{\Lambda_b}}[X_V^\lambda + Y_V^\lambda] , \\ b_V^\lambda &= -m_V g_V \frac{g_2}{m_{\Lambda_b}}[X_V^\lambda - Y_V^\lambda] , \\ x_V^\lambda &= m_V g_V [f_1 - \frac{m_{F_1} + m_{\Lambda_b}}{m_{\Lambda_b}} f_2][X_V^\lambda + Y_V^\lambda] , \\ y_V^\lambda &= -m_V g_V [g_1 + \frac{m_{\Lambda_b} - m_{F_1}}{m_{\Lambda_b}} g_2][X_V^\lambda - Y_V^\lambda] . \end{aligned} \tag{53}$$

Note that a_V^λ and x_V^λ now have the same weak phase, as do b_V^λ and y_V^λ.

Now consider again the triple-product terms of Eq. (36). As discussed above, to leading order in the heavy-quark expansion, only longitudinally-polarized vector mesons need be considered, and $\varepsilon_\mu^{\lambda=0} \sim q_\mu$ in this limit. Thus, we see that triple products of the form $\epsilon_{\alpha\beta\mu\nu} p_{F_1}^\alpha p_{\Lambda_b}^\beta s_{\Lambda_b}^\mu \varepsilon_V^\nu$ are of subleading order, and we neglect them. In fact, to leading order, there is only a single triple product which remains:

$$\begin{aligned} |\mathcal{M}_V|^2_{t.p.} \simeq & \frac{4}{m_V^2} \epsilon_{\alpha\beta\mu\nu} p_{\Lambda_b}^\alpha s_{\Lambda_b}^\beta q^\mu s_{F_1}^\nu \quad \left\{ -2\mathrm{Im}(ab^*)\, |q \cdot p_{\Lambda_b}|^2 + \mathrm{Im}\,(xy^*)\, q \cdot p_{\Lambda_b} \right. \\ & \left. + q \cdot p_{\Lambda_b} \left[(m_{\Lambda_b} + m_{F_1}) \mathrm{Im}\,(ay^*) + (m_{\Lambda_b} - m_{F_1}) \mathrm{Im}\,(bx^*) \right] \right\} . \end{aligned} \tag{54}$$

All other triple products are expected to be smaller, by a factor of order m_V/m_{Λ_b}.

We now turn to specific decays, and start with $\Lambda_b \to pK^{*-}$. First, for the tensor operators, one needs to evaluate matrix elements of the form

$$\langle V|\, \bar{q}_2 \sigma_{\mu\nu}(1 \pm \gamma_5) q_3 \,|0\rangle \,\langle F|\, \bar{q}_3 \sigma_{\mu\nu}(1 \pm \gamma_5) b \,|\Lambda_b\rangle \ . \tag{55}$$

However, as we have argued above, the tensor matrix element vanishes to leading order in the heavy-quark expansion. Therefore the tensor operators will not contribute to this decay. The same does not hold true for the scalar/pseudoscalar and vector/axial vector new-physics operators, and we find

$$\begin{aligned}
X_{K^*}^{NP,\lambda} &= \frac{G_F}{\sqrt{2}}\left[-\frac{1}{2}a_{u,1}^{LR} + b_{u,1}^{LL}\right], \\
Y_{K^*}^{\prime NP,\lambda} &= \frac{G_F}{\sqrt{2}}\left[-\frac{1}{2}a_{u,1}^{RL} + b_{u,1}^{RR}\right]. \qquad (56)
\end{aligned}$$

Note that, as expected, the new-physics operators contribute equally to longitudinally- and transversely-polarized V's. It is therefore reasonable to concentrate on the longitudinal V's, which dominate the decay $\Lambda_b \to pK^{*-}$.

The expressions for the decay $\Lambda_b \to n\bar{K}^{0*}$ can be easily obtained from those above by the replacement $u \to d$.

Finally, for $\Lambda_b \to \Lambda\phi$, we have

$$\begin{aligned}
X_{\phi}^{NP,\lambda} &= \frac{G_F}{\sqrt{2}}\left[-\frac{1}{2}a_{s,1}^{LR} + b_{s,1}^{LL} + b_{s,2}^{LL} + b_{s,2}^{LR}\right], \\
X_{\phi}^{SM,\lambda} &= -\frac{G_F}{\sqrt{2}}V_{tb}V_{ts}^*\left[a_3^t + a_4^t + a_5^t - \frac{1}{2}a_7^t - \frac{1}{2}a_9^t - \frac{1}{2}a_{10}^t\right. \\
&\qquad \left. - a_3^c - a_4^c - a_5^c + \frac{1}{2}a_7^c + \frac{1}{2}a_9^c + \frac{1}{2}a_{10}^c\right], \\
Y_{\phi}^{NP,\lambda} &= \frac{G_F}{\sqrt{2}}\left[-\frac{1}{2}a_{s,1}^{RL} + b_{s,1}^{RR} + b_{s,2}^{RR} + b_{s,2}^{RL}\right], \\
Y_{\phi}^{SM,\lambda} &\simeq 0\,, \qquad (57)
\end{aligned}$$

where we have included the standard model contribution without the tiny dipole contribution. The definitions of the various coefficients a_i^q, as well as their values, can be found in Ref. [5].

In all of the above decays, Y_V is expected to be very small in the SM, so that the triple products in $\Lambda_b \to F_1 V$ are predicted to be at most $O(1\%)$.

However, this can change significantly in the presence of new physics – from the above expressions one sees that the new-physics operators can easily produce a nonzero Y_V. The triple products in $\Lambda_b \to F_1 V$ may well be sizeable in the presence of new physics.

4 Diagnostic Power

In the previous section, we saw that the presence of new-physics operators can significantly modify the SM predictions for triple-product correlations in $b \to s$ Λ_b decays. In particular, triple products which were expected to be tiny in the SM may now be sizeable. This is not at all surprising: most of those triple products are vanishingly small because the decays are dominated by a single weak $b \to s$ penguin decay amplitude. However, in the presence of new physics, one can have several decay amplitudes and, consequently, large triple-product asymmetries.

Although this particular result is entirely expected, the previous exercise is still useful for several reasons. First, the pattern of nonzero triple products provides information about the type of new-physics operators which may be present. And second, one can apply the above general analysis to specific models of new physics to obtain model-dependent predictions. These are the issues we explore in this section.

We begin with the model-independent analysis. The first observation is simple: if one sees no new effect in a particular decay, this implies that certain new-physics operators are absent (barring fine-tuned cancellations among these operators). For example, suppose that the triple-product asymmetry in $\Lambda_b \to pK^{*-}$ is found to be tiny, as in the SM. This means that $Y_{K^*}^{NP} = 0$ [Eq. (56)], so that $a_{u,1}^{RL} = b_{u,1}^{RR} = 0$. (Note: since $Y_{K^*}^{SM} \simeq 0$, $X_{K^*}^{NP}$ could still be nonzero, since the triple product is proportional to the product of these two quantities.) This in turn suggests that each of $f_{u,1}^{RL}$, $f_{u,2}^{RL}$, $g_{u,1}^{RR}$ and $g_{u,2}^{RR}$ vanish, since they make up $a_{u,1}^{RL}$ and $b_{u,1}^{RR}$. Similarly, should no new effects be seen in $\Lambda_b \to \Lambda\eta$, each of the 30 operators in Y_η^{NP} [Eq. (45)] must vanish.

Of course, one can obtain more information by combining measurements, since the same operators can contribute to more than one decay. In fact, one can even partially test the assumption that there are no fine-tuned cancellations. For example, suppose that the triple-product asymmetry in $\Lambda_b \to pK^-$ is found to agree with the SM, but that in $\Lambda_b \to pK^{*-}$ does not. The latter result implies that $a_{u,1}^{RL}$ and/or $b_{u,1}^{RR}$ are nonzero. However, these operators also contribute to Y_K^{NP} [Eq. (41)]. Thus, in order to obtain $Y_K^{NP} = 0$, there must be cancellations among the various operators. Should such a result be found, it would be necessary to explain these cancellations, either via a symmetry, or by construction within a given model.

We now turn to the model-dependent analysis. The very general results of the previous section can be applied to specific models of new physics. Of course, in a given model, not all the operators of Eq. (2) will appear. In addition, it may be that the coefficients of those operators which do appear are related in some way. As examples of this behaviour, we examine those models described in Sec. 2, but this analysis can be applied to any models of new physics (e.g. supersymmetry, left-right symmetric models, etc.).

Consider first supersymmetric models with R-parity breaking (Sec. 2.1). If only B-violating couplings are present, then the only new-physics operators are vector operators contributing to $b \to s\bar{d}d$ [Eq. (8)]. This leads to a clear pattern of predictions: no new-physics effects are expected in the decays of a Λ_b to pK^-, pK^{*-} and $\Lambda\phi$. Indeed, if measurements of these triple-product asymmetries disagree with the SM predictions, this particular model is ruled out.

On the other hand, the decays $\Lambda_b \to \Lambda\eta$, $n\bar{K}^0$ and $n\bar{K}^{*0}$ can be affected in this model. How big can these effects be? In general, they can be enormous. As we have already noted, the new-physics contributions to these rare decays are still allowed by data to be comparable to, or even larger than, the SM contributions. If the two interfering amplitudes are of similar size, the triple-product asymmetry can be as large as $\sim 50\%$ (to be contrasted with the SM prediction of $\simeq 0$). This also holds for the other models discussed below.

Turning to the L-violating couplings, one sees that more operators may be present [Eqs. (14), (17), (20)]. In this case, all decays may be affected, except Λ_b to pK^-. This is a a quite distinctive signature for this model.

Finally, we consider leptophobic Z'-mediated FCNC's (Sec. 2.2). There are only six nonzero new-physics coefficients, given in Eq. (28), and these all depend on the parameters $\left|U_{sb}^{Z'}\right|$ and $M_{Z'}$ [Eq. (27)]. In this case, all Λ_b decays will be affected. However, note that, within this model, there are more observables (6) than there are theoretical parameters (2). This means that if deviations from the SM predictions are measured, we will be able to get a handle on $\left|U_{sb}^{Z'}\right|$ and $M_{Z'}$. Conversely, if no new-physics effects are observed, we will be able to place strong constraints on these quantities.

5 Conclusions

In the standard model (SM), (almost) all T-violating triple-product correlations in charmless Λ_b decays are expected to be tiny. (The one exception is the decay $\Lambda_b \to pK^-$, for which the asymmetry is 18%.) This is therefore a good place to look for physics beyond the standard model.

In this paper, using an effective-lagrangian approach, we have computed the effects of new physics on such triple products. This approach has the advantage of indicating which specific new-physics operators affect each of the Λ_b triple-product correlations. Thus, the measurement of a number of different triple products permits us to determine which new-physics operators are or are not present. Furthermore, the approach is completely general – the effects of any specific model can be obtained by simply calculating which operators appear in that model.

The new-physics effects on triple products are calculated using factorization. In addition, we work only to leading order in the heavy-quark expansion, neglecting terms of order m/m_{Λ_b}, where m is the mass of the light final-state meson. The justification for this is that it is only in fine-tuned scenarios that the new physics contributes at subleading order, but not at leading order. (Also, the subleading terms are quite a bit smaller, e.g. $m_{K^*}/m_{\Lambda_b} \sim 15\%$.)

We have found that all Λ_b triple products can be significantly modified by new physics. Of course, this to be expected. Most of the triple products are vanishingly small in the SM because the decays are dominated by a single weak $b \to s$ penguin decay amplitude. Thus, in the presence of new physics, there may be several decay amplitudes which can interfere with the SM amplitude. However, in order to obtain a sizeable asymmetry, the interfering amplitudes must be of similar size. We note that the constraints on the new-physics operators are sufficiently weak that they can be comparable to, or even larger than, the SM contributions. Thus, triple products which vanish in the SM can be as large as $\sim 50\%$ with new physics.

We have demonstrated how the measurement of triple-product asymmetries provides diagnostic information about the new-physics operators present. For example, all operators which affect $\Lambda_b \to pK^{*-}$ also affect $\Lambda_b \to pK^-$, but not vice-versa. Thus, if the triple product in $\Lambda_b \to pK^-$ is found to agree with the SM, we would also expect no new effects in $\Lambda_b \to pK^{*-}$. If this were found not to hold, then we would conclude that there must be cancellations among the operators in $\Lambda_b \to pK^-$, and this would have to be explained in some way (e.g. symmetry, specific model, etc.).

Finally, we have also applied this general approach to two specific models: supersymmetry with R-parity breaking, and leptophobic Z'-mediated flavour-changing neutral currents. In both cases, we have worked out the new-physics operators which appear in those models, and used the previous formalism to calculate which Λ_b triple products can be affected. For example, in the case of R-parity breaking models, there is a clear pattern of effects. One such model predicts significant new effects in the decays $\Lambda_b \to \Lambda\eta$, $n\bar{K}^0$ and $n\bar{K}^{*0}$, but not in Λ_b to pK^-, pK^{*-} and $\Lambda\phi$. Any deviation from this pattern would rule out this model. Other models of new physics can be treated similarly.

Acknowledgements: This work was financially supported by NSERC of Canada.

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CHAPITRE IV: QUATRIÈME ARTICLE

(Doit être publié dans: Physical Review D.)

Lepton Polarization and Forward-Backward Asymmetries in $b \to s\tau^+\tau^-$

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Abstract

We study the spin polarizations of both τ leptons in the decay $b \to s\tau^+\tau^-$. In addition to the polarization asymmetries involving a single τ, we construct asymmetries for the case where both polarizations are simultaneously measured. We also study forward-backward asymmetries with polarized τ's. We find that a large number of asymmetries are predicted to be large, $\gtrsim 10\%$. This permits the measurement of all Wilson coefficients and the b-quark mass, thus allowing the standard model (SM) to be exhaustively tested. Furthermore, there are many unique signals for the presence of new physics. For example, asymmetries involving triple-product correlations are predicted to be tiny within the SM, $O(10^{-2})$. Their observation would be a clear signal of new physics.

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1 Introduction

There has been a great deal of theoretical work examining the decay $b \to s\ell^+\ell^-$, both at the inclusive and exclusive level [1]. As usual, the hope is that, through precision measurements of this decay, one will find evidence for the presence of physics beyond the standard model (SM). Indeed, this decay mode has been extensively studied in various models of new physics [2].

Some years ago, it was noted that the measurement of the polarization of the final-state τ^- in the inclusive decay $b \to X_s\tau^+\tau^-$ can provide important information about the Wilson coefficients of the underlying effective Hamiltonian [3, 4, 5]. Within the SM, this inclusive decay is described in terms of five theoretical parameters: the four Wilson coefficients (C_7, C_{10} and real and imaginary parts of C_9), and the mass of the b-quark, m_b. In principle, all of these theoretical parameters can be completely determined using measurements of the three τ^- polarization asymmetries, the total (unpolarized) rate, and the forward-backward (FB) asymmetry.

In practice, however, the SM τ^- polarization asymmetry along the normal component is expected to be $O(10^{-2})$ [6], and is therefore probably too small to be measured. This situation can be remedied to some extent if, in addition to the polarization asymmetries of the τ^-, we also consider similar asymmetries for the τ^+ [7]. This adds one more independent observable. However, even if the sizeable polarization asymmetries of both τ^+ and τ^- can be separately measured, there are only as many measurements as there are unknowns, so that there are no redundant measurements to provide crosschecks for the SM. Furthermore, this program requires that the flavor of the b-quark be tagged: in an untagged sample, there are only four observables, since the measurement of the FB asymmetry requires tagging. It will therefore be very difficult to rigorously test the SM if only single τ-polarization measurements are made in

$b \to X_s \tau^+ \tau^-$.

In this paper, we try to construct the maximum possible number of independent observables. This is achieved by considering the situation in which both τ^+ and τ^- polarizations are simultaneously measured. As we will see, a variety of new asymmetries can be constructed in this case. We compute the polarization and forward-backward asymmetries for both singly-polarized and doubly-polarized final-state leptons. A large number of these new asymmetries do not require the tagging of the b-quark. (Note that, in an untagged sample, while the FB asymmetry for unpolarized leptons vanishes, some of the FB asymmetries for polarized leptons are nonvanishing.) On the other hand, if b-tagging is possible, the measurement of these new asymmetries provides even more information. The polarized FB asymmetries as well as the double-spin polarization asymmetries all depend in different ways on the Wilson coefficients, so that these coefficients can be obtained in many different ways. This redundancy provides a huge number of crosschecks, and allows the SM to be exhaustively tested. An interesting consequence of the large number of observables, is that m_b can be extracted. If the phenomenologically-obtained value of m_b were to agree with theoretical estimates [8], this would be an important step in confirming our understanding of QCD,

In our calculations we consider only contributions from SM operators. However, using arguments based on CPT invariance and the properties of the SM operators under C, P and T, we derive relations between these observables which are clean tests of new physics. Some of these tests rely on the fact that within the SM there are negligible CP-violating contributions to the decay mode being considered. Our philosophy is to test for the presence of new physics (NP) without considering the detailed structure of the various operators that can contribute to NP. Should a signal for NP be seen, the consideration of specific NP operators would help in determining the nature of

NP contributions (for example, see Ref. [7]).

We begin in Sec. 2 with a discussion of the calculation of $|\mathcal{M}|^2$, where $\mathcal{M}$ is the amplitude for $b \to s\tau^+\tau^-$ (the results of the calculation of $\mathcal{M}$ are complicated, and are presented in the Appendix). The polarization asymmetries and forward-backward asymmetries are examined in Secs. 3 and 4, respectively. We discuss these asymmetry measurements within a variety of scenarios in Sec. 5. We conclude in Sec. 6.

In total, we numerically evaluate the differential decay rate and 31 asymmetries as a function of the invariant lepton mass. Note that it will be extremely difficult to measure asymmetries smaller than 10%, as they would require $\sim 10^{10}$ B mesons for a 3σ signal (not including efficiencies for spin-polarization measurements and for tagging). We therefore consider only asymmetries larger than 10% as measurable. If one can only measure an individual τ^+ or τ^- spin, but cannot tag the flavor of the b-quark, then there are only two sizeable observables. If b-tagging can be done and one can measure the spin of the τ^+ or τ^-, this increases to 6 measurable asymmetries. However, if the polarizations of both the τ-leptons can be measured and flavor tagging of the b is possible, we find that nine of the asymmetries constructed here are large in the SM. Including the decay rate, this leads to 10 sizeable observables, which allows for a redundant test of the SM.

In addition, we find that the violation of certain SM asymmetry relations are clean tests of NP. Some of these relations are violated only in the presence of CP-violating NP. A large numbers of these asymmetries are $O(10^{-2})$ in the SM, so that the observation of larger asymmetries would be signals of NP. Certain combinations of these asymmetries are identically zero in the SM, and hence are litmus tests of NP.

2 $|\mathcal{M}|^2$ for $b \to s\tau^+\tau^-$

We begin by considering the calculation of $|\mathcal{M}|^2$ for $b \to s\tau^+\tau^-$. Including QCD corrections, the effective Hamiltonian describing the decay $b \to s\tau^+\tau^-$ [9] leads to the matrix element

$$\mathcal{M} = T_9 + T_{10} + T_7 \;, \tag{1}$$

where

$$T_9 = \frac{\alpha G_F}{\sqrt{2}\pi} C_9^{eff} V_{tb} V_{ts}^* \left[\bar{s}_L \gamma_\mu b_L\right] \left[\bar{\tau}^- \gamma^\mu \tau^+\right], \tag{2}$$

$$T_{10} = \frac{\alpha G_F}{\sqrt{2}\pi} C_{10} V_{tb} V_{ts}^* \left[\bar{s}_L \gamma_\mu b_L\right] \left[\bar{\tau}^- \gamma^\mu \gamma_5 \tau^+\right], \tag{3}$$

$$T_7 = \frac{\alpha G_F}{\sqrt{2}\pi} C_7^{eff} V_{tb} V_{ts}^* \left[\frac{-2im_b}{q^2}\right] \left[\bar{s}_L \sigma_{\mu\nu} q^\nu b_R\right] \left[\bar{\tau}^- \gamma^\mu \tau^+\right] . \tag{4}$$

In the above, q is the momentum transferred to the lepton pair, and we have neglected the s-quark mass. The Wilson coefficients C_i are evaluated perturbatively at the electroweak scale and then evolved down to the renormalization scale μ. The coefficients C_7^{eff} and C_{10} are real, and take the values

$$C_7^{eff} = -0.315 \;, \qquad C_{10} = -4.642 \tag{5}$$

in the leading-logarithm approximation [5]. On the other hand, the coefficient C_9^{eff} is complex, and its value is a function of $\hat{s} \equiv q^2/m_b^2$: $C_9^{eff}(\mu) \equiv C_9(\mu) + Y(\mu, \hat{s})$, where the function $Y(\mu, \hat{s})$ contains the one-loop contributions of the four-quark operators [3, 9]. An additional contribution to C_9^{eff} arises from the long-distance effects associated with real $c\bar{c}$ resonances in the intermediate states [10]. Thus, within the SM, the decay $b \to s\tau^+\tau^-$ is described by four Wilson coefficients for a given value of $\hat{s}$: C_7^{eff}, C_{10}, Re(C_9^{eff}) and Im(C_9^{eff}).

Because the expressions in $|\mathcal{M}|^2$ are complicated, we present the actual results of this calculation in the Appendix. Note that the polarization and forward-backward asymmetries, which will be discussed in subsequent sections, are calculated as functions of the terms of $|\mathcal{M}|^2$. Also, some signals of new physics are derived using the C, P and T properties of the terms at the $|\mathcal{M}|^2$ level.

There is one point which is worth mentioning here. In the calculation of $|\mathcal{M}|^2$, there are terms which involve the imaginary pieces of the Wilson coefficients [e.g. the $\mathrm{Im}(C_9^{eff} C_{10}^*)$ term in Eq. (51)]. These are the coefficients of terms like $\epsilon_{\mu\alpha\beta\phi}\, p_s^\mu p_-^\alpha s_-^\beta p_+^\phi$ in $|\mathcal{M}|^2$, which give rise to triple-product correlations (e.g. $\vec{p}_- \cdot (\vec{p}_+ \times \vec{s}_-)$). Naively, these triple products appear to violate time-reversal symmetry (T) and so, by the CPT theorem, should also be signals of CP violation. However, all the amplitudes in Eq. (1) have the same weak phase (neglecting the small u-quark contribution in the loop), so that their interference cannot give rise to CP violation. Thus, there appears to be an inconsistency.

What is happening is the following: a triple product is not a true T-violating signal, since the action of T exchanges the initial and final states. Because of this, triple-product correlations can be faked by the presence of strong phases, even if there is no CP violation. This is the situation which arises here – nonzero strong phases of the Wilson coefficients can lead to triple products. Usually, it is CP violation which interests us, and we wish to eliminate such fake signals. However, in this case, we are interested in measuring the imaginary parts of the Wilson coefficients in order to test the SM, so that these fake signals will be quite useful.

3 Polarization Asymmetries

In the computation of the various polarization asymmetries we choose a frame of reference in which the leptons move back to back along the z-axis, with the τ^- moving in the direction $+\hat{z}$. The s-quark then goes in the same direction as the b-quark, with the s-quark making an angle θ with the τ^-. Our specific choices for the 4-momenta components are as follows:

$$
\begin{aligned}
p^\mu_{\tau^-} &= \{\sqrt{P^2+m_\tau^2}, 0, 0, P\} , \\
p^\mu_{\tau^+} &= \{\sqrt{P^2+m_\tau^2}, 0, 0, -P\} , \\
p^\mu_{s} &= \{K, 0, K\sin\theta, K\cos\theta\} , \\
p^\mu_{b} &= \{\sqrt{K^2+m_b^2}, 0, K\sin\theta, K\cos\theta\}. \qquad (6)
\end{aligned}
$$

Using the above calculation of $|\mathcal{M}|^2$, we can compute the decay rate for unpolarized leptons by summing over the lepton spins and integrating over the angular variables. As a function of the invariant mass of the lepton pair, this decay rate is given by

$$
\left(\frac{d\Gamma(\hat{s})}{d\hat{s}}\right)_{\text{unpol}} = \frac{G_F^2\, m_b^5}{192\,\pi^3}\frac{\alpha^2}{4\,\pi^2}\,|V_{tb}\,V_{ts}^*|^2\,(1-\hat{s})^2\sqrt{1-\frac{4\,\hat{m}_\tau{}^2}{\hat{s}}}\,\Delta\ , \qquad (7)
$$

where $\hat{m}_\tau \equiv m_\tau/m_b$, and

$$
\begin{aligned}
\Delta &= \left(12\,\mathrm{Re}(C_7^{eff}C_9^{eff^*}) + \frac{4\,|C_7^{eff}|^2\,(2+\hat{s})}{\hat{s}}\right)\left(1+\frac{2\,\hat{m}_\tau^2}{\hat{s}}\right) \qquad (8) \\
&\quad +(|C_9^{eff}|^2+|C_{10}|^2)\left(1+2\,\hat{s}+\frac{2\,(1-\hat{s})\,\hat{m}_\tau^2}{\hat{s}}\right)+6\,(|C_9^{eff}|^2-|C_{10}|^2)\,\hat{m}_\tau^2.
\end{aligned}
$$

This agrees with the earlier results [3, 4, 5, 9, 11] in the appropriate limits.

We now consider the possibility that the polarizations of the final-state leptons can be measured. The spins of the $\tau^\pm$ are defined in their rest frames to be:

$$
\tilde{s}^\mu_{\tau^-} = \left\{0, s_x^-, s_y^-, s_z^-\right\}\ ,\quad \tilde{s}^\mu_{\tau^+} = \left\{0, s_x^+, s_y^+, s_z^+\right\}\ . \qquad (9)
$$

One can obtain the spins of the $\tau^\pm$ in the frame of Eq. (6) straightforwardly by performing a Lorentz boost:

$$
\begin{aligned}
s^{\mu}_{\tau^-} &= \left\{ \frac{P}{m_\tau} s_z^-, s_x^-, s_y^-, \frac{\sqrt{P^2+m_\tau^2}}{m_\tau} s_z^- \right\} , \\
s^{\mu}_{\tau^+} &= \left\{ -\frac{P}{m_\tau} s_z^+, s_x^+, s_y^+, \frac{\sqrt{P^2+m_\tau^2}}{m_\tau} s_z^+ \right\} . \qquad (10)
\end{aligned}
$$

We now define differential decay rate as a function of the spin directions of the $\tau^\pm$, $\mathbf{s}^+$ and $\mathbf{s}^-$, where $\mathbf{s}^+$ and $\mathbf{s}^-$ are unit vectors in the $\tau^\pm$ rest frames. This is given by

$$
\begin{aligned}
\frac{d\Gamma(\mathbf{s}^+, \mathbf{s}^-, \hat{s})}{d\hat{s}} &= \frac{1}{4}\left(\frac{d\Gamma(\hat{s})}{d\hat{s}}\right)_{\text{unpol}} \Bigg[1 + \Big(\mathcal{P}_x^- \, s_x^- + \mathcal{P}_y^- \, s_y^- + \mathcal{P}_z^- \, s_z^- \\
&\quad + \mathcal{P}_x^+ \, s_x^+ + \mathcal{P}_y^+ \, s_y^+ + \mathcal{P}_z^+ \, s_z^+\Big) \\
&\quad + \Big(\mathcal{P}_{xx} \, s_x^+ s_x^- + \mathcal{P}_{xy} \, s_x^+ s_y^- + \mathcal{P}_{xz} \, s_x^+ s_z^- + \mathcal{P}_{yx} \, s_y^+ s_x^- \\
&\quad + \mathcal{P}_{yy} \, s_y^+ s_y^- + \mathcal{P}_{yz} \, s_y^+ s_z^- + \mathcal{P}_{zx} \, s_z^+ s_x^- + \mathcal{P}_{zy} \, s_z^+ s_y^- + \mathcal{P}_{zz} \, s_z^+ s_z^-\Big)\Bigg], \qquad (11)
\end{aligned}
$$

where the single-lepton polarization asymmetries $\mathcal{P}_i^\mp (i = x, y, z)$ are obtained by evaluating

$$
\begin{aligned}
\mathcal{P}_i^- &= \frac{\left[\frac{d\Gamma(\mathbf{s}^-=\hat{\mathbf{i}},\, \mathbf{s}^+=\hat{\mathbf{i}})}{d\hat{s}} + \frac{d\Gamma(\mathbf{s}^-=\hat{\mathbf{i}},\, \mathbf{s}^+=-\hat{\mathbf{i}})}{d\hat{s}}\right] - \left[\frac{d\Gamma(\mathbf{s}^-=-\hat{\mathbf{i}},\, \mathbf{s}^+=\hat{\mathbf{i}})}{d\hat{s}} + \frac{d\Gamma(\mathbf{s}^-=-\hat{\mathbf{i}},\, \mathbf{s}^+=-\hat{\mathbf{i}})}{d\hat{s}}\right]}{\left[\frac{d\Gamma(\mathbf{s}^-=\hat{\mathbf{i}},\, \mathbf{s}^+=\hat{\mathbf{i}})}{d\hat{s}} + \frac{d\Gamma(\mathbf{s}^-=\hat{\mathbf{i}},\, \mathbf{s}^+=-\hat{\mathbf{i}})}{d\hat{s}}\right] + \left[\frac{d\Gamma(\mathbf{s}^-=-\hat{\mathbf{i}},\, \mathbf{s}^+=\hat{\mathbf{i}})}{d\hat{s}} + \frac{d\Gamma(\mathbf{s}^-=-\hat{\mathbf{i}},\, \mathbf{s}^+=-\hat{\mathbf{i}})}{d\hat{s}}\right]} , \\
\mathcal{P}_i^+ &= \frac{\left[\frac{d\Gamma(\mathbf{s}^-=\hat{\mathbf{i}},\, \mathbf{s}^+=\hat{\mathbf{i}})}{d\hat{s}} + \frac{d\Gamma(\mathbf{s}^-=-\hat{\mathbf{i}},\, \mathbf{s}^+=\hat{\mathbf{i}})}{d\hat{s}}\right] - \left[\frac{d\Gamma(\mathbf{s}^-=\hat{\mathbf{i}},\, \mathbf{s}^+=-\hat{\mathbf{i}})}{d\hat{s}} + \frac{d\Gamma(\mathbf{s}^-=-\hat{\mathbf{i}},\, \mathbf{s}^+=-\hat{\mathbf{i}})}{d\hat{s}}\right]}{\left[\frac{d\Gamma(\mathbf{s}^-=\hat{\mathbf{i}},\, \mathbf{s}^+=\hat{\mathbf{i}})}{d\hat{s}} + \frac{d\Gamma(\mathbf{s}^-=-\hat{\mathbf{i}},\, \mathbf{s}^+=\hat{\mathbf{i}})}{d\hat{s}}\right] + \left[\frac{d\Gamma(\mathbf{s}^-=\hat{\mathbf{i}},\, \mathbf{s}^+=-\hat{\mathbf{i}})}{d\hat{s}} + \frac{d\Gamma(\mathbf{s}^-=-\hat{\mathbf{i}},\, \mathbf{s}^+=-\hat{\mathbf{i}})}{d\hat{s}}\right]} . \qquad (12)
\end{aligned}
$$

Similarly, the double spin asymmetries $\mathcal{P}_{ij}$ can be obtained:

$$
\mathcal{P}_{ij} = \frac{\left[\frac{d\Gamma(\mathbf{s}^+=\hat{\mathbf{i}},\mathbf{s}^-=\hat{\mathbf{j}})}{d\hat{s}} - \frac{d\Gamma(\mathbf{s}^+=\hat{\mathbf{i}},\mathbf{s}^-=-\hat{\mathbf{j}})}{d\hat{s}}\right] - \left[\frac{d\Gamma(\mathbf{s}^+=-\hat{\mathbf{i}},\mathbf{s}^-=\hat{\mathbf{j}})}{d\hat{s}} - \frac{d\Gamma(\mathbf{s}^+=-\hat{\mathbf{i}},\mathbf{s}^-=-\hat{\mathbf{j}})}{d\hat{s}}\right]}{\left[\frac{d\Gamma(\mathbf{s}^+=\hat{\mathbf{i}},\mathbf{s}^-=\hat{\mathbf{j}})}{d\hat{s}} + \frac{d\Gamma(\mathbf{s}^+=\hat{\mathbf{i}},\mathbf{s}^-=-\hat{\mathbf{j}})}{d\hat{s}}\right] + \left[\frac{d\Gamma(\mathbf{s}^+=-\hat{\mathbf{i}},\mathbf{s}^-=\hat{\mathbf{j}})}{d\hat{s}} + \frac{d\Gamma(\mathbf{s}^+=-\hat{\mathbf{i}},\mathbf{s}^-=-\hat{\mathbf{j}})}{d\hat{s}}\right]} , \qquad (13)
$$

where $\hat{i}$ and $\hat{j}$ are unit vectors along the i and j directions. Note that both $\mathcal{P}_i^\pm$ and $\mathcal{P}_{ij}$ depend also on $\hat{s}$. However, the explicit dependence on $\hat{s}$ has been suppressed for simplicity of notation.

Before presenting explicit expressions for these quantities, it is useful to make the following remark. With our choice of 4-momenta [Eq. (6)], the decay takes place in the yz plane. Therefore, the only vectors which can have $\hat{x}$ components are the spins $\mathbf{s}^+$ and $\mathbf{s}^-$. This implies that the only scalar product which involves $\hat{x}$ components is the dot product of two spins. Thus, any term that has only one component of spin along $\hat{x}$ (i.e. $\mathcal{P}_x$, $\mathcal{P}_{xy}$ and $\mathcal{P}_{xz}$) must come from a triple-product correlation. This holds even in the presence of new physics. It is therefore these quantities which probe the imaginary parts of the products of Wilson coefficients.

The $\mathcal{P}$'s take the form

$$\mathcal{P}_x^+ = \frac{-3\pi}{2\sqrt{\hat{s}}\Delta}\Big(2\,\mathrm{Im}(C_7^{eff}C_{10}^*) + \mathrm{Im}(C_9^{eff}C_{10}^*)\,\hat{s}\Big)\,\hat{m}_\tau\sqrt{1-\frac{4\,\hat{m}_\tau^2}{\hat{s}}} \tag{14}$$

$$\begin{aligned}\mathcal{P}_y^+ = \frac{3\pi}{2\sqrt{\hat{s}}\Delta}\Big(&\frac{4\,|C_7^{eff}|^2}{\hat{s}} + 2\,\mathrm{Re}(C_7^{eff}C_{10}^*) + 4\,\mathrm{Re}(C_7^{eff}C_9^{eff^*}) \\ &+ \mathrm{Re}(C_9^{eff}C_{10}^*) + |C_9^{eff}|^2\,\hat{s}\Big)\,\hat{m}_\tau\end{aligned} \tag{15}$$

$$\mathcal{P}_z^+ = \frac{2}{\Delta}\Big(6\,\mathrm{Re}(C_7^{eff}C_{10}^*) + \mathrm{Re}(C_9^{eff}C_{10}^*)\,(1+2\,\hat{s})\Big)\sqrt{1-\frac{4\,\hat{m}_\tau^2}{\hat{s}}} \tag{16}$$

$$\mathcal{P}_x^- = \mathcal{P}_x^+ \tag{17}$$

$$\begin{aligned}\mathcal{P}_y^- = \frac{3\pi}{2\sqrt{\hat{s}}\Delta}\Big(&\frac{4\,|C_7^{eff}|^2}{\hat{s}} - 2\,\mathrm{Re}(C_7^{eff}C_{10}^*) + 4\,\mathrm{Re}(C_7^{eff}C_9^{eff^*}) \\ &- \mathrm{Re}(C_9^{eff}C_{10}^*) + |C_9^{eff}|^2\,\hat{s}\Big)\,\hat{m}_\tau\end{aligned} \tag{18}$$

$$\mathcal{P}_z^- = \mathcal{P}_z^+ \tag{19}$$

$$\mathcal{P}_{xx} = \frac{1}{\Delta}\Bigg(24\,\mathrm{Re}(C_7^{eff}C_9^{eff^*})\,\frac{\hat{m}_\tau^2}{\hat{s}} + 4\,|C_7^{eff}|^2\,\frac{\big((-1+\hat{s})\,\hat{s} + 2\,(2+\hat{s})\,\hat{m}_\tau^2\big)}{\hat{s}^2}$$

$$+(|C_9^{eff}|^2-|C_{10}|^2)\,\frac{\left((1-\hat{s})\hat{s}+2\,(1+2\,\hat{s})\hat{m}_\tau^2\right)}{\hat{s}}\Bigg) \tag{20}$$

$$\mathcal{P}_{yx} = \frac{-2}{\Delta}\,\mathrm{Im}(C_9^{eff}C_{10}^*)\,(1-\hat{s})\sqrt{1-\frac{4\,\hat{m}_\tau^2}{\hat{s}}} \tag{21}$$

$$\mathcal{P}_{zx} = \frac{-3\,\pi}{2\,\sqrt{\hat{s}}\Delta}\left(2\,\mathrm{Im}(C_7^{eff}C_{10}^*)+\mathrm{Im}(C_9^{eff}C_{10}^*)\right)\hat{m}_\tau \tag{22}$$

$$\mathcal{P}_{xy} = \mathcal{P}_{yx} \tag{23}$$

$$\begin{aligned}\mathcal{P}_{yy} = \frac{1}{\Delta}\Bigg(&\frac{24\,\mathrm{Re}(C_7^{eff}C_9^{eff^*})\,\hat{m}_\tau^2}{\hat{s}}-4\,(|C_9^{eff}|^2+|C_{10}|^2)\,\frac{(1-\hat{s})\,\hat{m}_\tau^2}{\hat{s}}\\ &+(|C_9^{eff}|^2-|C_{10}|^2)\,((-1+\hat{s})+\frac{6\,\hat{m}_\tau^2}{\hat{s}})\\ &+\frac{4\,|C_7^{eff}|^2\left((1-\hat{s})\,\hat{s}+2\,(2+\hat{s})\,\hat{m}_\tau^2\right)}{\hat{s}^2}\Bigg)\end{aligned} \tag{24}$$

$$\begin{aligned}\mathcal{P}_{zy} = \frac{3\,\pi}{2\,\sqrt{\hat{s}}\Delta}&\left(2\,\mathrm{Re}(C_7^{eff}C_{10}^*)-|C_{10}|^2+\mathrm{Re}(C_9^{eff}C_{10}^*)\,\hat{s}\right)\\ &\times\hat{m}_\tau\sqrt{1-\frac{4\,\hat{m}_\tau^2}{\hat{s}}}\end{aligned} \tag{25}$$

$$\mathcal{P}_{xz} = -\mathcal{P}_{zx} \tag{26}$$

$$\begin{aligned}\mathcal{P}_{yz} = \frac{3\,\pi}{2\,\sqrt{\hat{s}}\Delta}&\left(2\,\mathrm{Re}(C_7^{eff}C_{10}^*)+|C_{10}|^2+\mathrm{Re}(C_9^{eff}C_{10}^*)\,\hat{s}\right)\\ &\times\hat{m}_\tau\sqrt{1-\frac{4\,\hat{m}_\tau^2}{\hat{s}}}\end{aligned} \tag{27}$$

$$\begin{aligned}\mathcal{P}_{zz} = \frac{1}{\Delta}\Bigg(&12\,\mathrm{Re}(C_7^{eff}C_9^{eff^*})\left(1-\frac{2\,\hat{m}_\tau^2}{\hat{s}}\right)+\frac{4\,|C_7^{eff}|^2\,(2+\hat{s})\left(1-\frac{2\hat{m}_\tau^2}{\hat{s}}\right)}{\hat{s}}\\ &+(|C_9^{eff}|^2+|C_{10}|^2)\left(1+2\,\hat{s}-\frac{6\,(1+\hat{s})\,\hat{m}_\tau^2}{\hat{s}}\right)\\ &+\frac{2\,(|C_9^{eff}|^2-|C_{10}|^2)\,(2+\hat{s})\,\hat{m}_\tau^2}{\hat{s}}\Bigg)\,.\end{aligned} \tag{28}$$

The coefficient $\mathcal{P}_z^-$ was computed in Ref. [4], $\mathcal{P}_x^-$, $\mathcal{P}_y^-$ and $\mathcal{P}_z^-$ were obtained in Ref. [5], and $\mathcal{P}_x^+$, $\mathcal{P}_y^+$ and $\mathcal{P}_z^+$ were calculated in Ref. [7]. (Note: while we agree with the calculations of Refs. [4, 5], we disagree with Ref. [7] about

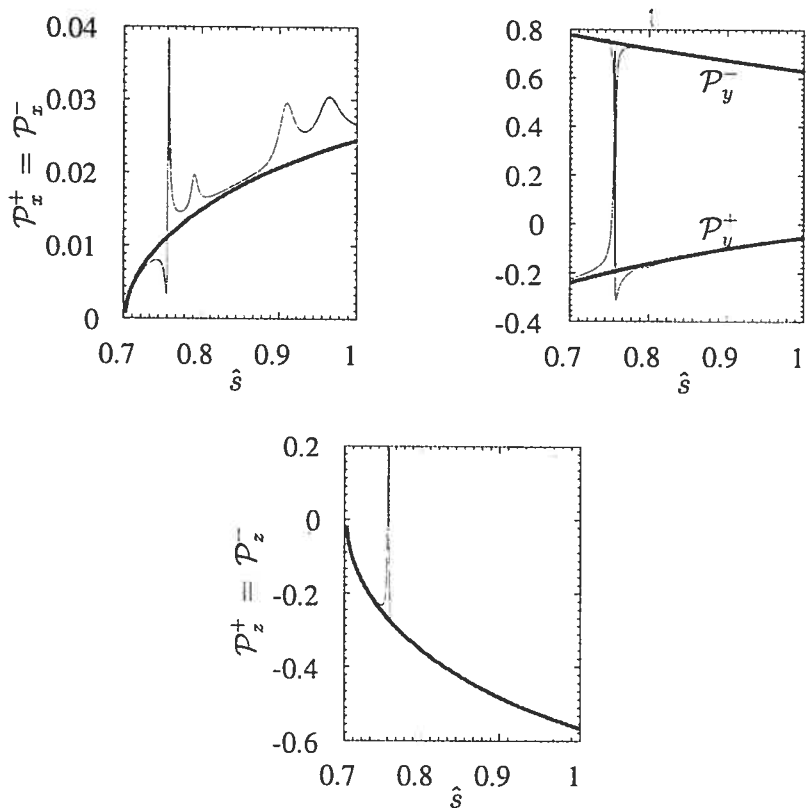

Figure 1: The polarization asymmetries for the τ^- and τ^+, as functions of $\hat{s}$, the invariant mass of the τ pair, without (thick lines) and with (thin lines) the long-distance resonance contributions.

the expression for $\mathcal{P}_y^{\pm}$ [the equation following their Eq. (24)].) We plot the functions $\mathcal{P}_x^-$, $\mathcal{P}_y^-$ and $\mathcal{P}_z^-$ as functions of $\hat{s}$ in Fig. 1. For the purpose of numerical computations, we follow the prescription of Ref. [5] and include long-distance effects in C_9^{eff} associated with real $c\bar{c}$ resonances in the intermediate states. We take the phenomenological parameter κ_V multiplying the Breit-Wigner function in Ref. [5] to be unity.

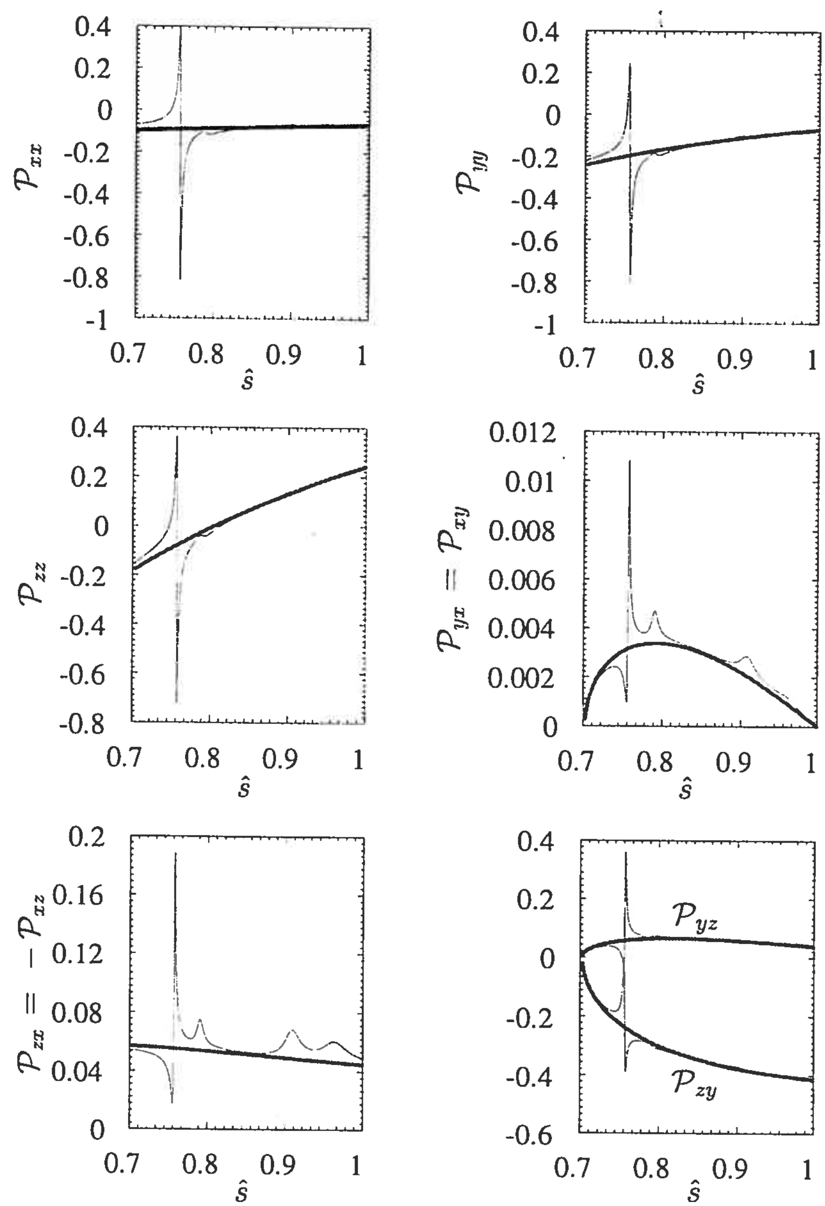

Figure 2: The double-spin polarization asymmetries, as functions of $\hat{s}$, the invariant mass of the τ pair, without (thick lines) and with (thin lines) the long-distance resonance contributions.

Note that our $\mathcal{P}_z^-$ is the same as the longitudinal polarization asymmetry of the τ^-, P_L^- of Refs. [5, 7]. However, $\mathcal{P}_z^+ = -P_L^+$, since the τ^+ moves along the $-\hat{z}$ axis. Similarly, $\mathcal{P}_x^- = P_N^-$ and $\mathcal{P}_x^+ = -P_N^+$. (Note: the distribution of our P_N differs from that of Ref. [5], resulting in a somewhat smaller value of $\langle P_N \rangle_{\tau}$.) The transverse direction defined in these references lies along the negative $\hat{y}$-direction for both τ^- and τ^+, so that $\mathcal{P}_y^- = -P_T^-$ and $\mathcal{P}_y^+ = -P_T^+$. The double-spin polarization asymmetries are shown in Fig. 2.

Previously, we noted that it is very likely that only asymmetries larger than 10% will be measurable. One characterization of the data is to calculate the average values of the above asymmetries[5]. These are defined as

$$\langle \mathcal{P} \rangle \equiv \frac{\displaystyle\int_{4\hat{m}_\tau^2}^{1} \mathcal{P} \frac{d\Gamma}{d\hat{s}} d\hat{s}}{\displaystyle\int_{4\hat{m}_\tau^2}^{1} \frac{d\Gamma}{d\hat{s}} d\hat{s}} \, . \tag{29}$$

In Table 1 we list the average values of all polarization asymmetries. From this Table, we see that only $\mathcal{P}_y^{\pm}$, $\mathcal{P}_z$, $\mathcal{P}_{yy}$, and $\mathcal{P}_{zy}$ can be considered sizeable. (Note that here, $\mathcal{P}_z^+ = \mathcal{P}_z^- \equiv \mathcal{P}_z$.)

In Eqs. (14)–(28), there are certain relations between the $\mathcal{P}$'s when the spins s^+ and s^- are interchanged. Some of these relations are equalities, e.g. $\mathcal{P}_x^- = \mathcal{P}_x^+$, $\mathcal{P}_{xz} = \mathcal{P}_{zx}$, etc. For other pairs of $\mathcal{P}_i$'s, the expressions are similar, but only some of the terms change sign (e.g. $\mathcal{P}_y^+$ vs. $\mathcal{P}_y^-$). As we describe below, it is possible to understand these relations by considering also the conjugate process $\bar{b} \to \bar{s}\tau^-\tau^+$.

The processes $b \to s\tau^+\tau^-$ and $\bar{b} \to \bar{s}\tau^-\tau^+$ are related by CPT as follows [12]:

$$b(p_b) \;\to\; s(p_s)\, \tau^+(p_+, \mathrm{s}_+)\, \tau^-(p_-, \mathrm{s}_-) \; ,$$

[5]It is also possible that the average value of an asymmetry is small, but that large values of the asymmetry are still possible for certain values of $\hat{s}$. For example, see $\mathcal{P}_{zz}$.

$\langle \mathcal{P}_x^- \rangle = \langle \mathcal{P}_x^+ \rangle$	1.413×10^{-2}
$\langle \mathcal{P}_y^- \rangle$	0.723
$\langle \mathcal{P}_z^- \rangle = \langle \mathcal{P}_z^+ \rangle$	-0.336
$\langle \mathcal{P}_y^+ \rangle$	-0.164
$\langle \mathcal{P}_{xx} \rangle$	-8.658×10^{-2}
$\langle \mathcal{P}_{yx} \rangle = \langle \mathcal{P}_{xy} \rangle$	2.868×10^{-3}
$\langle \mathcal{P}_{zx} \rangle = -\langle \mathcal{P}_{xz} \rangle$	5.322×10^{-2}
$\langle \mathcal{P}_{yy} \rangle$	-0.168
$\langle \mathcal{P}_{zy} \rangle$	-0.281
$\langle \mathcal{P}_{yz} \rangle$	5.717×10^{-2}
$\langle \mathcal{P}_{zz} \rangle$	-1.1254×10^{-2}

Table 1: Numerical values of the various averaged spin-polarization asymmetries without including the long-distance resonance contributions. We use $m_b = 4.24$ GeV [8]. The corresponding branching ratio is $BR(B \to X_s \tau^+ \tau^-) = 1.192 \times 10^{-7}$.

$$\bar{b}(p_b) \;\to\; \bar{s}(p_s)\, \tau^-(p_+, -\mathrm{s}_+)\, \tau^+(p_-, -\mathrm{s}_-) \; . \tag{30}$$

In the absence of CP violation, observables which are P-odd must vanish in the (C-even) untagged sample. Consider first the terms involving triple-product (TP) correlations. While all triple products are T-odd, they can be either P-even or P-odd. Triple products involving two spins are necessarily P-odd and, in the absence of CP violation, C-odd. Because of this, in the SM, these triple product must vanish in the untagged sample. Thus, we have $\mathrm{TP}_{\bar{b}}^{P-odd} = -\mathrm{TP}_{b}^{P-odd}$. This relation can be violated in the presence of CP-violating new physics. On the other hand, triple products involving one spin are P-even and C-even, so that $\mathrm{TP}_{\bar{b}}^{P-even} = +\mathrm{TP}_{b}^{P-even}$, in the absence of CP violation. Thus, these triple products can survive in the untagged sample due to the presence of the strong phases which can fake CP-violating effects.

We now apply these observations to $\mathcal{P}_x^+$ and $\mathcal{P}_x^-$, which involve a single spin. As noted earlier, terms with a single spin along $\hat{x}$ must come only from a triple-product correlation. The general triple-product term giving these quantities can be written as $\epsilon_{\alpha\beta\mu\rho}\, p_b^\alpha p_s^\beta (a\, p_+^\mu s_+^\rho + b\, p_-^\mu s_-^\rho)$, where a and b are arbitrary coefficients. For the conjugate process [Eq. (30)], the corresponding term is $-\epsilon_{\alpha\beta\mu\rho}\, p_{\bar{b}}^\alpha p_+^\beta (a\, p_-^\mu s_-^\rho + b\, p_+^\mu s_+^\rho)$. Since $\mathrm{TP}_{\bar{b}}^{P-even} = +\mathrm{TP}_{b}^{P-even}$, this implies that $a = -b$ (in the absence of CP violation). Using the 4-vectors of Eq. (6), it is then straightforward to show that this results in $\mathcal{P}_x^+ = +\mathcal{P}_x^-$. Note that this will hold even in the presence of CP-conserving New Physics.

Similarly, the two-spin triple products, which contribute to the pairs $\{\mathcal{P}_{yx}, \mathcal{P}_{xy}\}$ and $\{\mathcal{P}_{zx}, \mathcal{P}_{xz}\}$, are proportional to $\epsilon_{\alpha\beta\mu\rho}\, p_s^\alpha p_b^\beta s_-^\mu s_+^\rho$. In the absence of CP violation, the CP-odd combination of $\mathcal{P}_{yx}$ and $\mathcal{P}_{xy}$ (and of $\mathcal{P}_{zx}$ and $\mathcal{P}_{xz}$) will vanish in an untagged sample. Again, a simple calculation then shows that this implies that $\mathcal{P}_{yx} = +\mathcal{P}_{xy}$ and $\mathcal{P}_{zx} = -\mathcal{P}_{xz}$.

For the other terms that do not contain triple products, and are hence always T-even, one can understand the relationship between the $\mathcal{P}$'s in a similar fashion. For example, consider $\mathcal{P}_y^+$ and $\mathcal{P}_y^-$. Since only dot products of various momenta and one spin are involved, the coefficients of both terms $|C_7^{eff}|^2$ [Eq. (52)] and $\mathrm{Re}(C_7^{eff} C_{10}^*)$ [Eq. (54)] are T-even and P-odd. However, the $|C_7^{eff}|^2$ term "$p_s \cdot (s^- + s^+)$" switches sign under CPT for the conjugate process, while the $\mathrm{Re}(C_7^{eff} C_{10}^*)$ term "$p_s \cdot (s^+ - s^-)$" has the same sign for the conjugate process. Since these terms are P-odd and C-odd (in the absence of CP violation), they must vanish in an untagged sample. This explains the relative sign difference between the $|C_7^{eff}|^2$ and $\mathrm{Re}(C_7^{eff} C_{10}^*)$ terms in $\mathcal{P}_y^+$ and

$\mathcal{P}_y^-$. This argument may be extended to all terms contributing to various $\mathcal{P}_i$'s. In particular, in the SM, $\mathcal{P}_z^+ = +\mathcal{P}_z^-$. On the other hand, in presence of New Physics, while the additional terms must still be T-even and P-odd, they could be even or odd under CPT, implying that the relation between $\mathcal{P}_z^+$ and $\mathcal{P}_z^-$ could differ.

Of course, the above discussion assumes that there is no CP violation in $b \to s\tau^+\tau^-$, which is the case in the SM, to a good approximation. On the other hand, if new CP-violating physics contributes to this decay, this gives us several clear tests for its presence. For example, any violation of the relation $\mathcal{P}_x^+ = \mathcal{P}_x^-$ (or $P_L^- + P_L^+ = 0$) is a smoking-gun signal of such new physics.

4 Forward-Backward Asymmetries

One observable which does not depend on the polarization of the final-state leptons is the forward-backward (FB) asymmetry. In the frame of reference described in Eq. (6), the forward-backward asymmetry is given by

$$\begin{aligned} A_{FB}(\hat{s}) &= \frac{\displaystyle\int_0^1 \frac{d^2\Gamma}{d\hat{s}\; d\cos\theta} d\cos\theta - \int_{-1}^0 \frac{d^2\Gamma}{d\hat{s}\; d\cos\theta} d\cos\theta}{\displaystyle\int_0^1 \frac{d^2\Gamma}{d\hat{s}\; d\cos\theta} d\cos\theta + \int_{-1}^0 \frac{d^2\Gamma}{d\hat{s}\; d\cos\theta} d\cos\theta} \\ &= \frac{3}{\Delta}\left(2\,\mathrm{Re}(C_7^{eff} C_{10}^*) + \hat{s}\,\mathrm{Re}(C_9^{eff} C_{10}^*)\right)\sqrt{1 - \frac{4\,\hat{m}_\tau^2}{\hat{s}}}\,. \end{aligned} \tag{31}$$

This agrees with the result of Ref. [5] (and that of Ref. [11] when m_τ is neglected). Note that the FB asymmetry is of opposite sign for the CP-conjugate process $\bar{b} \to \bar{s}\tau^+\tau^-$, so that $A_{FB}^b + A_{FB}^{\bar{b}} = 0$. Thus, in order to measure the unpolarized FB asymmetry, it will be necessary to tag the flavor of the decaying b-quark.

If the polarization of the final-state leptons can be measured, then, in addition to the polarization asymmetries discussed in the previous section, one can also extract forward-backward asymmetries of the polarized leptons. While the unpolarized FB asymmetry of Eq. (31) requires b-tagging, some of the polarized FB asymmetries are non-vanishing even in an untagged sample.

We can extract the forward-backward asymmetries corresponding to various polarization components of the τ^- and/or τ^+ spin by writing:

$$\begin{aligned} A_{FB}(\mathbf{s}^+, \mathbf{s}^-, \hat{s}) &= A_{FB}(\hat{s}) + \Big[\mathcal{A}_x^- s_x^- + \mathcal{A}_y^- s_y^- + \mathcal{A}_z^- s_z^- \\ &\quad + \mathcal{A}_x^+ s_x^+ + \mathcal{A}_y^+ s_y^+ + \mathcal{A}_z^+ s_z^+ \\ &\quad + \mathcal{A}_{xx} s_x^+ s_x^- + \mathcal{A}_{xy} s_x^+ s_y^- + \mathcal{A}_{xz} s_x^+ s_z^- \\ &\quad + \mathcal{A}_{yx} s_y^+ s_x^- + \mathcal{A}_{yy} s_y^+ s_y^- + \mathcal{A}_{yz} s_y^+ s_z^- \\ &\quad + \mathcal{A}_{zx} s_z^+ s_x^- + \mathcal{A}_{zy} s_z^+ s_y^- + \mathcal{A}_{zz} s_z^+ s_z^- \Big] \, . \end{aligned} \qquad (32)$$

The various polarized forward-backward asymmetries are then evaluated to be

$$\mathcal{A}_x^+ = 0 \qquad (33)$$

$$\mathcal{A}_y^+ = \frac{2}{\Delta} \mathrm{Re}(C_9^{eff} C_{10}^*) \frac{(1-\hat{s})\, \hat{m}_\tau}{\sqrt{\hat{s}}} \sqrt{1 - \frac{4\,\hat{m}_\tau^2}{\hat{s}}} \qquad (34)$$

$$\begin{aligned} \mathcal{A}_z^+ &= \frac{1}{\Delta}\Big(6\,\mathrm{Re}(C_7^{eff} C_9^{eff*}) - \frac{6\,|C_7^{eff}|^2}{\hat{s}} - 3\,(\,|C_9^{eff}|^2 - |C_{10}|^2)\,\hat{m}_\tau^2 \\ &\quad -12\,\mathrm{Re}(C_7^{eff} C_{10}^*)\,\frac{\hat{m}_\tau^2}{\hat{s}} - 6\,\mathrm{Re}(C_9^{eff} C_{10}^*)\,\frac{\hat{m}_\tau^2}{\hat{s}} \\ &\quad - \frac{3}{2}\,(\,|C_9^{eff}|^2 + |C_{10}|^2)\,\hat{s}\,(1 - \frac{2\,\hat{m}_\tau^2}{\hat{s}})\Big) \end{aligned} \qquad (35)$$

$$\mathcal{A}_x^- = 0 \qquad (36)$$

$$\mathcal{A}_y^- = \mathcal{A}_y^+ \qquad (37)$$

$$\begin{aligned} \mathcal{A}_z^- &= \frac{1}{\Delta}\Big(- 6\,\mathrm{Re}(C_7^{eff} C_9^{eff*}) - \frac{6\,|C_7^{eff}|^2}{\hat{s}} - 3\,(\,|C_9^{eff}|^2 - |C_{10}|^2)\,\hat{m}_\tau^2 \\ &\quad + 12\,\mathrm{Re}(C_7^{eff} C_{10}^*)\,\frac{\hat{m}_\tau^2}{\hat{s}} + 6\,\mathrm{Re}(C_9^{eff} C_{10}^*)\,\frac{\hat{m}_\tau^2}{\hat{s}} \end{aligned}$$

$$- \frac{3}{2} \left(|C_9^{eff}|^2 + |C_{10}|^2 \right) \hat{s} \left(1 - \frac{2\,\hat{m}_\tau^2}{\hat{s}}\right) \Big) \tag{38}$$

$$\mathcal{A}_{xx} = 0 \tag{39}$$

$$\mathcal{A}_{xy} = \frac{-6}{\Delta} \left(2\,\mathrm{Im}(C_7^{eff}\, C_{10}^*) + \mathrm{Im}(C_9^{eff}\, C_{10}^*)\right) \frac{\hat{m}_\tau^2}{\hat{s}} \tag{40}$$

$$\mathcal{A}_{xz} = \frac{2}{\Delta}\,\mathrm{Im}(C_9^{eff} C_{10}^*)\, \frac{(1-\hat{s})\,\hat{m}_\tau}{\sqrt{\hat{s}}} \sqrt{1 - \frac{4\,\hat{m}_\tau^2}{\hat{s}}} \tag{41}$$

$$\mathcal{A}_{yx} = -\mathcal{A}_{xy} \tag{42}$$

$$\mathcal{A}_{yy} = 0 \tag{43}$$

$$\mathcal{A}_{yz} = \left(2|C_9^{eff}|^2 - \frac{8\,|C_7^{eff}|^2}{\hat{s}}\right) \frac{(1-\hat{s})\,\hat{m}_\tau}{\Delta\sqrt{\hat{s}}} \tag{44}$$

$$\mathcal{A}_{zx} = \mathcal{A}_{xz} \tag{45}$$

$$\mathcal{A}_{zy} = \mathcal{A}_{yz} \tag{46}$$

$$\mathcal{A}_{zz} = \frac{-3}{\Delta} \left(2\,\mathrm{Re}(C_7^{eff}\, C_{10}^*) + \mathrm{Re}(C_9^{eff}\, C_{10}^*)\,\hat{s}\right) \sqrt{1 - \frac{4\,\hat{m}_\tau^2}{\hat{s}}}\,. \tag{47}$$

Note that, in the SM, it turns out that $A_{FB} = -\mathcal{A}_{zz}$.

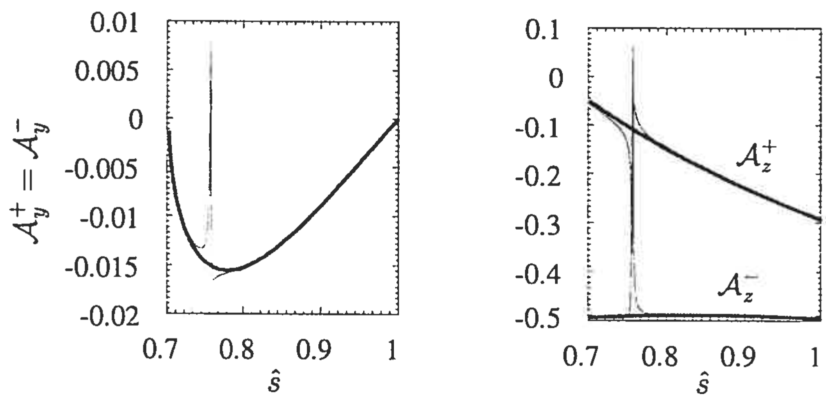

Figure 3: Forward-backward asymmetries of the τ^- and τ^+, as functions of $\hat{s}$, the invariant mass of the τ pair, without (thick lines) and with (thin lines) the long-distance resonance contributions.

The nonzero single-spin forward-backward asymmetries are depicted in

Fig. 3 as functions of $\hat{s}$, while those with both spins polarized are shown in Fig. 4. Interestingly, some of the forward-backward asymmetries are identically zero within the SM: $\mathcal{A}_x^+$, $\mathcal{A}_x^-$, $\mathcal{A}_{xx}$ and $\mathcal{A}_{yy}$. Nonvanishing values of these asymmetries would be clear signals of NP. Also, as was discussed in the case of the polarization asymmetries, the discrete transformation properties of the operators can once again be used to understand the relations between pairs of forward-backward asymmetries in which the spins s^+ and s^- are interchanged.

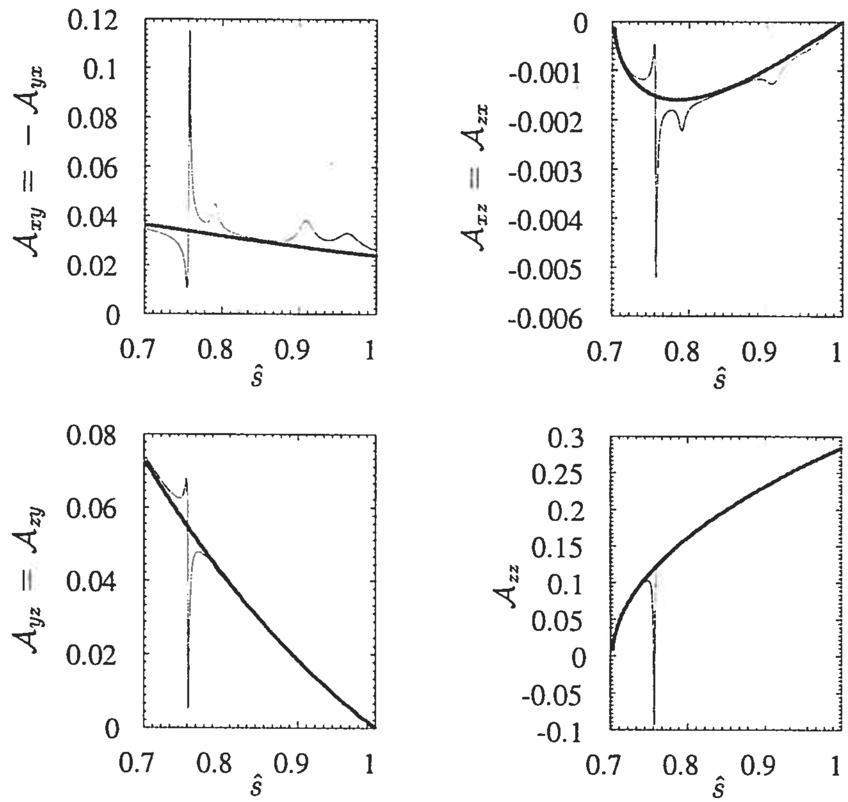

Figure 4: Doubly-polarized forward-backward asymmetries, as functions of $\hat{s}$, the invariant mass of the τ pair, without (thick lines) and with (thin lines) the long-distance resonance contributions.

The average values of the forward-backward asymmetries are defined similarly to Eq. (29) and are listed in Table 2. From this table, we see that only three asymmetries are expected to be larger than 10% in the SM: $\mathcal{A}_z^{\pm}$ and $\mathcal{A}_{zz}$.

$\langle \mathcal{A}_y^+ \rangle = \langle \mathcal{A}_y^- \rangle$	-1.302×10^{-2}
$\langle \mathcal{A}_z^+ \rangle$	-0.148
$\langle \mathcal{A}_z^- \rangle$	-0.490
$\langle \mathcal{A}_{xy} \rangle = -\langle \mathcal{A}_{yx} \rangle$	3.184×10^{-2}
$\langle \mathcal{A}_{xz} \rangle = \langle \mathcal{A}_{zx} \rangle$	-1.347×10^{-3}
$\langle \mathcal{A}_{yz} \rangle = \langle \mathcal{A}_{zy} \rangle$	4.298×10^{-2}
$\langle \mathcal{A}_{zz} \rangle$	0.154

Table 2: Numerical values of the various average polarized forward-backward asymmetries without including the long distance resonance contributions. We use $m_b = 4.24$ GeV [8]. The corresponding average unpolarized forward-backward asymmetry is $\langle A_{FB} \rangle = -0.154$.

5 Discussion

In the previous sections, we have discussed the polarization and forward-backward asymmetries which can be obtained when the spins of the τ^+ and/or τ^- are measured. Here we consider what can be learned from these measurements in a variety of scenarios.

First, suppose that the statistics are such that only a single polarization can be measured (say that of the τ^-), and that no tagging is possible. In this case only the P-even observables survive: $\mathcal{P}_x^+ + \mathcal{P}_x^-$, $\mathcal{A}_y^+ + \mathcal{A}_y^-$ and $\mathcal{A}_z^+ + \mathcal{A}_z^-$.

Of these asymmetries only $\mathcal{A}_z^+ + \mathcal{A}_z^-$ is measurable within the SM. Along with the differential decay rate, this therefore gives only two observables, which is not enough to test the SM.

On the other hand, if the polarizations of both τ^+ and τ^- can be measured, still without tagging, then one adds another six observables: $\mathcal{P}_{xx}$, $\mathcal{P}_{yy}$, $\mathcal{P}_{zy} + \mathcal{P}_{yz}$, $\mathcal{P}_{zz}$, $\mathcal{A}_{xy} + \mathcal{A}_{yx}$, and $\mathcal{A}_{xz} + \mathcal{A}_{zx}$. Of these only three — $\mathcal{P}_{yy}$, $\mathcal{P}_{zy} + \mathcal{P}_{yz}$ and $\mathcal{P}_{zz}$ — are expected to be sizeable in the SM, the last one being measurable only as a distribution in $\hat{s}$ (see Fig. 2). We therefore have just enough measurements to determine the five unknowns C_7^{eff}, C_{10}, $\mathrm{Re}(C_9^{eff})$, $\mathrm{Im}(C_9^{eff})$, and m_b. However, there are not enough measurements to provide an internal crosscheck of the predictions of the SM.

Now suppose that it is possible to tag the flavor of the decaying b-quark. If only a single τ-spin measurement is performed then, out of a total of thirteen possible asymmetries, only six are sizeable within the SM: $\mathcal{A}_{FB}$, $\mathcal{P}_y^{\mp}$, $\mathcal{P}_z$ and $\mathcal{A}_z^{\pm}$. (Recall that $\mathcal{P}_z^+ = \mathcal{P}_z^- \equiv \mathcal{P}_z$.)

In the best-case scenario, it will be possible to both tag the flavor of the decaying b, and to measure the polarizations of both final-state τ leptons. In this case, one in principle has 31 asymmetries. However, within the SM only nine of these are accessible: $\mathcal{A}_{FB}$, $\mathcal{P}_y^{\pm}$, $\mathcal{P}_z$, $\mathcal{A}_z^{\pm}$, $\mathcal{P}_{yy}$, $\mathcal{P}_{zy}$ and $\mathcal{A}_{zz}$. Even so, if these asymmetries could be measured, this would allow us to greatly overconstrain the SM. Ideally, we will find evidence for new physics, but if not, these will provide precision determinations of both m_b and the Wilson coefficients describing the decay $b \to s\tau^+\tau^-$.

In Table 3 we summarize the number of possible observables, including the differential cross-section, in the various scenarios discussed above.

Finally, we note that in some of these scenarios, it will be possible to extract the value of m_b. This is advantageous for two reasons. First, it permits a

		Untagged Sample	Tagged Sample
Only one of τ^+ or τ^- spin measured		$\frac{d\Gamma}{d\hat{s}}$, $\mathcal{P}_x^{(\pm)}$, $\mathcal{A}_y^{(\pm)}$, $\mathcal{A}_z^{(\pm)}$[4]	$\frac{d\Gamma}{d\hat{s}}$, $\mathcal{A}_{FB}$, $\mathcal{P}_x^{\pm}$, $\mathcal{P}_y^{\pm}$, $\mathcal{P}_z^{\pm}$, $\mathcal{A}_x^{\pm}$, $\mathcal{A}_y^{\pm}$, $\mathcal{A}_z^{\pm}$ [14]
	SM	$\frac{d\Gamma}{d\hat{s}}$, $\mathcal{A}_z^{(\pm)}$ [2]	$\frac{d\Gamma}{d\hat{s}}$, $\mathcal{A}_{FB}$, $\mathcal{P}_y^{\pm}$, $\mathcal{P}_z$, $\mathcal{A}_z^{\pm}$ [7]
Both τ^+ and τ^- spins measured		+ $\mathcal{P}_{xx}$, $\mathcal{P}_{yy}$, $\mathcal{P}_{(zy)}$, $\mathcal{P}_{zz}$, $\mathcal{A}_{(xy)}$, $\mathcal{A}_{(xz)}$ [10]	All [32]
	SM	+ $\mathcal{P}_{yy}$, $\mathcal{P}_{(zy)}$, $\mathcal{P}_{zz}$ [5]	+ $\mathcal{P}_{yy}$, $\mathcal{P}_{zy}$, $\mathcal{P}_{zz}$ [10]

Table 3: The number of observables in various scenarios of initial b-flavor tagging and τ-spin measurements. The columns represent the cases with untagged and tagged samples, while the rows are for the scenarios in which only one of the τ^+ or τ^- spin is measured, or when both τ^+ and τ^- spins are measured. Of the possible observables, those that are sizeable in the SM are listed separately. In the case in which both spins are measured, only the additional observables (indicated by +) are listed. The number in the square brackets represents the total number of observables possible in each case. $\mathcal{P}_{(ij)}$ indicates the sum $\mathcal{P}_{ij} + \mathcal{P}_{ji}$ and $\mathcal{P}_i^{(\pm)} = \mathcal{P}_i^+ + \mathcal{P}_i^-$, with identical definitions for the $\mathcal{A}$'s.

direct comparison with the theoretical estimates of m_b [8]. Second, for some measurements in the B system, it is necessary to input m_b from theory, which increases the systematic (theoretical) uncertainty of the measurement. By contrast, we see that the double-spin analysis of $b \to s\tau^+\tau^-$ will not suffer from this type of systematic error.

6 Conclusions

In the standard model (SM), the inclusive decay $b \to s\tau^+\tau^-$ is described by five theoretical parameters: C_7^{eff}, C_{10}, $\mathrm{Re}(C_9^{eff})(\hat{s})$, $\mathrm{Im}(C_9^{eff})(\hat{s})$ and m_b, where $\hat{s}$ is related to the momentum transferred to the lepton pair. We would like to be able to test this description.

In this paper, we have calculated all single- and double-spin asymmetries in the decay $b \to s\tau^+\tau^-$. We have shown that there are many different ways of testing the SM description of this decay. In all, there are a total of 31 different asymmetries. However, only 9 of these are predicted to be measurable, i.e. have values larger than 10%. (Indeed some asymmetries are expected to vanish in the SM.) Should any of the small asymmetries be found to have large values, this would be a clear signal of new physics (NP). Furthermore, the SM predicts certain relationships among the asymmetries when the spins s^+ and s^- are interchanged. Should these relations be violated, this would also indicate the presence of NP. In fact, this could give us some clue as to whether the NP is CP-conserving or CP-violating.

Apart from these signals of NP, whether or not the SM can be tested depends crucially on which types of measurements can be made. For example, if one cannot perform b-tagging, and can measure only a single individual τ spin, then there are only two sizeable observables. This is not enough to test the SM. On the other hand, if one can measure both τ spins, but cannot tag the flavor of the b, then there are a total of five measurable observables. This is enough to determine the theoretical unknowns, but does not provide the necessary redundancy to test the SM.

On the other hand, if one can perform b-tagging, but can only measure a single τ spin, then there are 7 sizeable observables. This can provide a redundant test of the SM. The optimal scenario is if b-tagging is possible, and one can measure the polarizations of both the τ^+ and τ^-. In this case, there are a total of 10 independent measurements, which would greatly overconstrain the SM. If new physics is not found, this would precisely determine the five theoretical parameters.

Note that testing the SM implies that the quantity m_b will be extracted from the experimental data. This will allow us to compare the experimental value of m_b with the theoretical estimates of this same quantity. Furthermore, as the measurements do not rely on theoretical input, the systematic error will be correspondingly reduced.

Acknowledgements: N.S. and R.S. thank D.L. for the hospitality of the Université de Montréal, where part of this work was done. The work of D.L. was financially supported by NSERC of Canada. The work of Nita Sinha was supported by a project of the Department of Science and Technology, India, under the young scientist scheme.

7 Appendix

In this Appendix, we calculate the square of the amplitude in Eq. (1), keeping the spins of both final-state leptons. We define p_b, p_s, p_+ and p_- to be the momenta of the b-quark, s-quark, τ^+ and τ^-, respectively, with $q = p_b - p_s = p_+ + p_-$. The spins of the τ^+ and τ^- are denoted by s_+ and s_-, respectively. We have

$$|\mathcal{M}|^2 = |T_9|^2 + |T_{10}|^2 + |T_7|^2 + 2\mathrm{Re}\left(T_9^\dagger T_{10}\right) + 2\mathrm{Re}\left(T_9^\dagger T_7\right) + 2\mathrm{Re}\left(T_{10}^\dagger T_7\right) \ . \quad (48)$$

Summing over the s-quark spin and averaging over the b-quark spin, We find

$$\frac{1}{2}\sum_{b,s\ spins}|T_9|^2 = \frac{\alpha^2 G_F^2}{\pi^2}|V_{tb}V_{ts}^*|^2|C_9^{eff}|^2$$
$$\times\left\{\frac{(m_b^2-q^2)}{2}\left(-p_-\cdot s_+\,p_+\cdot s_- + \frac{q^2}{2}s_+\cdot s_- + m_\tau^2(1-s_+\cdot s_-)\right)\right.$$
$$+(1-s_+\cdot s_-)\left(p_b\cdot p_+\,p_s\cdot p_- + p_s\cdot p_+\,p_b\cdot p_-\right)$$
$$-\frac{q^2}{2}\left[p_b\cdot s_+\,p_s\cdot s_- + p_s\cdot s_+\,p_b\cdot s_-\right]$$
$$+s_+\cdot p_-\left[p_b\cdot p_+\,p_s\cdot s_- + p_s\cdot p_+\,p_b\cdot s_-\right]$$
$$+s_-\cdot p_+\left[p_b\cdot p_-\,p_s\cdot s_+ + p_s\cdot p_-\,p_b\cdot s_+\right]$$
$$+m_\tau\Big[p_s\cdot(p_+ + p_-)\,p_b\cdot(s_+ + s_-)$$
$$\left.-p_b\cdot(p_+ + p_-)\,p_s\cdot(s_+ + s_-)\Big]\right\}, \qquad (49)$$

$$\frac{1}{2}\sum_{b,s\ spins}|T_{10}|^2 = \frac{\alpha^2 G_F^2}{\pi^2}|V_{tb}V_{ts}^*|^2|C_{10}|^2$$
$$\times\left\{-\frac{(m_b^2-q^2)}{2}\left(-p_-\cdot s_+\,p_+\cdot s_- + \frac{q^2}{2}s_+\cdot s_- + m_\tau^2(1-s_+\cdot s_-)\right)\right.$$
$$+(1+s_+\cdot s_-)\left(p_b\cdot p_+\,p_s\cdot p_- + p_s\cdot p_+\,p_b\cdot p_-\right)$$
$$-\left(2m_\tau^2-\frac{q^2}{2}\right)\left[p_b\cdot s_+\,p_s\cdot s_- + p_s\cdot s_+\,p_b\cdot s_-\right]$$
$$-s_+\cdot p_-\left[p_b\cdot p_+\,p_s\cdot s_- + p_s\cdot p_+\,p_b\cdot s_-\right]$$
$$-s_-\cdot p_+\left[p_b\cdot p_-\,p_s\cdot s_+ + p_s\cdot p_-\,p_b\cdot s_+\right]$$
$$-m_\tau\Big[p_s\cdot(p_+ - p_-)\,p_b\cdot(s_+ - s_-)$$
$$\left.-p_b\cdot(p_+ - p_-)\,p_s\cdot(s_+ - s_-)\Big]\right\}, \qquad (50)$$

$$\sum_{b,s\ spins} \mathrm{Re}\Big[T_9^\dagger T_{10}\Big] = \frac{\alpha^2 G_F^2}{\pi^2}|V_{tb}V_{ts}^*|^2$$
$$\times \Big\{2\mathrm{Re}(C_9^{eff}C_{10}^*)\Big[m_\tau^2[p_s\cdot s_-\,p_b\cdot s_+ - p_b\cdot s_-\,p_s\cdot s_+]$$
$$-\frac{q^2}{2}(p_s\cdot p_- - p_s\cdot p_+)$$
$$+m_\tau[p_b\cdot p_+\,p_s\cdot s_- + p_s\cdot p_+\,p_b\cdot s_-$$
$$-p_b\cdot p_-\,p_s\cdot s_+ - p_s\cdot p_-\,p_b\cdot s_+]\Big]$$
$$+\mathrm{Im}(C_9^{eff}C_{10}^*)\epsilon_{\mu\alpha\beta\phi}\Big[\Big[2m_\tau + (p_s+p_b)\cdot s_+\Big]p_s^\mu p_-^\alpha s_-^\beta p_+^\phi$$
$$-\Big[2m_\tau + (p_s+p_b)\cdot s_-\Big]p_s^\mu p_+^\alpha s_+^\beta p_-^\phi$$
$$-(p_s+p_b)\cdot p_+\,p_s^\mu p_-^\alpha s_-^\beta s_+^\phi + (p_s+p_b)\cdot p_-\,p_s^\mu p_+^\alpha s_+^\beta s_-^\phi$$
$$+(p_- - p_+)\cdot p_s\,p_-^\mu p_+^\alpha s_+^\beta s_-^\phi\Big]\Big\}\,, \qquad (51)$$

$$\frac{1}{2}\sum_{b,s\ spins}|T_7|^2 = \frac{\alpha^2 G_F^2}{\pi^2}|V_{tb}V_{ts}^*|^2\frac{m_b^2}{q^4}|C_7|^2$$
$$\times\Big\{4m_b^2 m_\tau\Big[p_s\cdot(p_+ + p_-)\,q\cdot(s_- + s_+) - q^2\,p_s\cdot(s_- + s_+)\Big]$$
$$+2m_\tau^2 m_b^2(m_b^2 - q^2)[(1 - s_+\cdot s_-] + q^2 m_b^2(m_b^2 - q^2)$$
$$-2q^2\Big[2[p_s\cdot p_-\,p_b\cdot p_+ + p_s\cdot p_+\,p_b\cdot p_-] - q^2\,p_b\cdot p_s\Big][(1 - s_+\cdot s_-]$$
$$-4q^2\Big[s_+\cdot p_-[p_s\cdot s_-\,p_b\cdot p_+ + p_b\cdot s_-\,p_s\cdot p_+]$$
$$+s_-\cdot p_+[p_s\cdot p_-\,p_b\cdot s_+ + p_b\cdot p_-\,p_s\cdot s_+]\Big]$$
$$+2q^2(m_b^2 - q^2)\,s_+\cdot p_-\,s_-\cdot p_+$$
$$+2q^4[p_s\cdot s_-\,p_b\cdot s_+ + p_b\cdot s_-\,p_s\cdot s_+]\Big\}\,, \qquad (52)$$

$$\sum_{b,s\ spins} \mathrm{Re}\Big[T_9^\dagger T_7\Big] = \frac{\alpha^2 G_F^2}{\pi^2}|V_{tb}V_{ts}^*|^2 \frac{m_b^2}{q^2}$$
$$\times \Big\{\mathrm{Re}(C_9^{eff} c_7^*)\Big[(m_b^2 - q^2)[q^2 + 2m_\tau^2 - 2m_\tau^2\, s_+ \cdot s_-]$$
$$-4m_\tau[p_b \cdot (p_+ + p_-)\, p_s \cdot (s_+ + s_-)$$
$$-p_s \cdot (p_+ + p_-)\, p_b \cdot (s_+ + s_-)]\Big]\Big\}\,, \qquad (53)$$

$$\sum_{b,s\ spins} \mathrm{Re}\Big[T_{10}^\dagger T_7\Big] = \frac{\alpha^2 G_F^2}{\pi^2}|V_{tb}V_{ts}^*|^2 \frac{m_b^2}{q^2}$$
$$\Big\{-2\mathrm{Re}(C_{10} C_7^{eff*})\Big[-\frac{m_\tau(m_b^2 - q^2)}{2}[p_+ \cdot s_- - p_- \cdot s_+]$$
$$+q^2[p_s \cdot p_- - p_s \cdot p_+] - 2m_\tau^2[p_s \cdot s_-\, p_- \cdot s_+ - p_+ \cdot s_-\, p_s \cdot s_+]$$
$$+m_\tau\, q^2 p_s \cdot (s_+ - s_-) + m_\tau\, p_s \cdot (p_- - p_+)(s_- \cdot p_+ + s_+ \cdot p_-)\Big]$$
$$-4\mathrm{Im}(C_{10} C_7^{eff*})\epsilon_{\alpha\beta\mu\rho}\Big[-m_\tau\, p_s^\alpha p_+^\beta s_+^\mu p_b^\rho + m_\tau\, p_s^\alpha p_-^\beta s_-^\mu p_b^\rho$$
$$+m_\tau^2\, s_-^\alpha p_s^\beta s_+^\mu p_b^\rho\Big]\Big\}\,. \qquad (54)$$

In the above, we have used the convention $\epsilon_{0123} = +1$. Note that, as mentioned in Sec. 2, C_7^{eff} and C_{10} are expected to be real; only C_9^{eff} is complex. However, in the expressions above, for completeness we have included both real and imaginary pieces of all Wilson coefficients.

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CONCLUSION

Les quatre articles de cette thèse renferment des résultats très intéressants dans la recherche sur la brisure CP/T. On s'est intéressé à une classe d'asymétries qui n'a pas été très exploitée et qui est fort prometteuse. C'est celle d'asymétries dues aux produits triples. On s'est penché sur les désintégrations des mésons B et des hypérons non-charmés Λ_b.

Dans le premier article on a considéré, dans le modèle standard (MS), le mode inclusif $b \to su\bar{u}$ avec tous les quarks polarisés. L'amplitude de désintégration a trois termes dominants: l'un provient d'un diagramme arbre ("*tree diagram*") et les deux autres d'un diagramme pingouin. On a calculé l'amplitude au carré en négligeant les masses finales, et on a trouvé que les termes d'interférences contiennent plusieurs produits triples formés des vecteurs spins et impulsions des quatre quarks. On ne s'est pas intéressé aux PT avec trois spins car trop difficiles à mesurer. Voici nos résultats principaux:

- De tous les PT trouvés, deux types seulement donnent des asymétries appréciables: ce sont $\vec{p}_u \cdot (\vec{s}_u \times \vec{s}_{\bar{u}})$ (ou $\vec{p}_{\bar{u}} \cdot (\vec{s}_u \times \vec{s}_{\bar{u}})$) et $\vec{s}_b \cdot (\vec{p}_u \times \vec{p}_s)$.
- Pour le premier type, l'asymétrie peut atteindre les 5%. On peut la rechercher dans les désintégrations $B \to V_1 V_2$ où V_1 et V_2 sont des bosons vecteurs. Ces modes ont été considérés dans la référence [6] de l'article et on peut dire que nos résultats sont cohérents avec les leurs.
- Pour le deuxième type, l'asymétrie peut atteindre les 3%. On ne peut pas la rechercher dans les désintégrations des mésons B car ces derniers sont sans spin et le spin du quark b intervient dans le PT. Par contre on peut rechercher ce signal dans les désintégrations des Λ_b, par exemple $\Lambda_b \to \Lambda\pi^+\pi^-$ ou $\Lambda_b \to \Lambda K^+ K^-$.

- Tous les termes des PT contenant le spin du quark s sont supprimés par un facteur d'au moins (m_s/m_b). Donc si un signal de PT contenant $\vec{s}_s$ est détecté dans les expériences ce serait une preuve de l'existence d'une nouvelle physique. On peut rechercher ce type de signaux dans le mode $\Lambda_b \to \Lambda\pi^+\pi^-$, via un PT faisant intervenir le spin du Λ puisque ce dernier vaut grossièrement celui du quark s.

- Les physiciens ne sont pas encore éclairés sur la façon dont les spins et impulsions d'un hadron sont reliés à ceux des quarks qui le constituent. En comparant nos résultats (niveau des quarks) aux résultats expérimentaux (niveau hadronique) on peut distinguer lequel des modèles de quarks est le plus proche de la réalité.

Dans le deuxième article, on a cherché les asymétries de brisure T dues aux PT, dans les modes $\Lambda_b \to FP$ et $\Lambda_b \to FV$ où F est un baryon léger, P un méson pseudoscalaire et V un méson vecteur. On a utilisé le modèle de factorisation en appliquant l'hamiltonien efficace des désintégrations des B non-charmés dans le MS à des modes spécifiques de désintégrations de Λ_b. On a tenu compte des effets à longue distance et on a utilisé l'approximation des quarks lourds ("*heavy-quark effective theory*"). L'étude aboutit aux résultats suivants:

- Dans les désintégrations $\Lambda_b \to FP$, il y a un seul produit triple dans l'amplitude au carré. Dans le référentiel propre de Λ_b, il s'écrit $\vec{p}_F \cdot (\vec{s}_F \times \vec{s}_{\Lambda_b})$, où $\vec{p}_F$ et $\vec{s}_F$ sont les vecteurs impulsion et spin de F et $\vec{s}_{\Lambda_b}$ est le vecteur spin de Λ_b. On a considéré trois modes de cette classe de désintégrations:

    1. Le mode $\Lambda_b \to pK^-$ (qui correspond au mode inclusif $b \to s u\bar{u}$) est le plus prometteur dans le MS. L'asymétrie due au PT $\vec{p}_p \cdot (\vec{s}_p \times \vec{s}_{\Lambda_b})$

peut atteindre les 18%. Ce PT peut être grossièrement identifié à $\vec{p}_u \cdot (\vec{s}_u \times \vec{s}_b)$ du mode inclusif. Il n'a pas été considéré dans le premier article car il fait partie des termes comprenant un facteur de suppression en $(1/m_b)$. Par contre, au niveau hadronique cette suppression est compensée par le terme d'accroissement chiral ("*the chiral enhancement*") $m_K^2/(m_s + m_u)$.

2. pour le mode $\Lambda_b \to p\pi^-$, l'étude est la même en remplaçant le quark s par le quark d. On trouve une asymétrie d'environ 5%.

3. On a aussi considéré les désintégrations $\Lambda_b \to \Lambda\eta(\eta')$ qui correspondent aux modes inclusifs $b \to sq\bar{q}$; où $q = u, d, s$. Ici l'asymétrie ne dépasse pas 1%. C'est un bon endroit pour chercher la nouvelle physique.

- Dans les désintégrations $\Lambda_b \to FV$, il y a plusieurs produits triples dans le carré de l'amplitude dont un similaire à celui de $\Lambda_b \to FP$. À part ce dernier, tous les autres PT contiennent le vecteur polarisation du boson V. Là aussi on a considéré trois modes spécifiques:

  1. Pour le mode $\Lambda_b \to pK^{*-}$ les asymétries sont de l'ordre de 1% pour K^{*-} polarisé transversalement et elles sont plus petites encore pour les K^{*-} longitudinaux.

  2. Les asymétries sont inférieures à 1% pour le mode $\Lambda_b \to p\rho^-$.

  3. Pour le processus $\Lambda_b \to \Lambda\phi$, dominé par une amplitude unique (un pûr pingouin), les asymétries sont nulles.

On voit que pour tous les modes $\Lambda_b \to FV$ étudiés, les asymétries de brisure T sont soit petites soit négligeables. C'est donc aussi un bon endroit pour chercher la nouvelle physique.

- On considère que l'utilisation du modèle de factorisation dans nos calculs est une bonne approximation car on a estimé que les contributions non-factorisables sont négligeables.

Dans le troisième article, on utilise la même méthode et on considère les mêmes modes que dans le deuxième article mais avec un lagrangien efficace contenant tous les opérateurs possibles de la nouvelle physique. On trouve que les nouveaux opérateurs peuvent changer considérablement les prédictions du MS, surtout pour les modes dominés (dans le MS) par une seule amplitude faible provenant d'un pingouin $b \to s$. La nouvelle physique leur apporte d'autres amplitudes et de là de grandes asymétries pouvant atteindre 50%.

Cette méthode a de grands avantages. Elle permet, en mesurant les differents signaux de PT, de décider quel opérateur de la nouvelle physique doit être présent dans l'hamiltonien et lequel doit être supprimé et ce, indépendamment de tout modèle. Ensuite, on peut appliquer nos résultats à n'importe quel modèle spécifique. Nous en avons pris deux pour exemples:

- Le premier modèle est celui de la supersymétrie avec brisure de la parité R. Si on ne considère que les couplages brisant le nombre baryonique, aucun effet au-delà du MS n'est prévu pour les désintégrations de Λ_b en pK^-, pK^{*-} et $\Lambda\phi$. Donc si les expériences donnent des asymétries dans ces modes différentes de celles prédites par le MS, le modèle en question devra être éliminé. Maintenant si on ne considère que les couplages brisant le nombre leptonique, tous les modes peuvent être affectés sauf $\Lambda_b \to pK^-$. Des mesures d'asymétries différentes de celles prédites par le MS dans ce mode élimineraient ce modèle.
- Le deuxième modèle est celui des FCNC's avec le Z' leptophobique comme médiateur. On a trouvé que ce modèle affectera tous les modes

de désintégration étudiés.

Dans le quatrième article, on a considéré la désintégration $b \to s\tau^+\tau^-$ avec les deux leptons polarisés et on a étudié, dans le MS, les asymétries de polarisation et "avant-arrière" ("*forward-backward*"). Les résultats de cette étude permettent de tester de différentes manières le MS et de sonder la nouvelle physique. En effet, il y a un total de 31 asymétries dans ce mode, dont 9 sont mesurables (de l'ordre de 10% et plus). En outre, il y a cinq paramètres théoriques qui décrivent le processus dans le MS: quatre coefficients de Wilson et la masse du quark b.

- Si on arrive a effectuer l'étiquetage ("*the tagging*") des quarks b, on aura la possibilité de mesurer 7 observables indépendantes dans le cas d'un seul lepton polarisé et 10 dans le cas des deux leptons polarisés. Cela donne assez de redondance pour tester les cinq paramètres du MS, ce qui permettra soit de trouver de la nouvelle physique, soit de donner des valeurs plus précises à ces paramètres. Ajoutons que la mesure de la masse du quark b réduira les erreurs systématiques dans les expériences sur les mésons B.

- Nous avons trouvé que certaines asymétries sont très petites et d'autres sont carrément nulles dans le MS. Donc si on trouve dans les expériences des asymétries plus grandes que les valeurs prédites, ce sera un signal de la présence d'une nouvelle physique. Un exemple intéressant est l'existence d'asymétries dues à des produits triples (ce sont de faux signaux de brisure CP car dûs seulement aux phases fortes) qui sont négligeables dans le MS (de l'ordre de 10^{-2}). Leur observation constituerait un signal évident de la nouvelle physique.

- Finalement, on a trouvé qu'il y avait certaines relations entre les asymétries lorsqu'on permute les spins des deux leptons. Si ces relations ne sont pas respectées en expérience, ce sera aussi une manifestation de la nouvelle physique.

Pour conclure, on peut dire que l'étude des asymétries CP dues aux produits triples offre plusieurs manières de tester le MS et de rechercher la nouvelle physique. Nous avons montré que cette étude pourrait donner de l'information sur l'origine de la brisure CP, inconnue jusqu'à présent.

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"Et dis: Ô mon Seigneur, accroîs mes connaissances "

[Saint Coran 20:114]