Grand Unified Theories and Beyond

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Standard Model very successful

→ low energy accessible part of a (more) Fundamental Theory of Elementary Particles.

BUT with

ad hoc Higgs sector
ad hoc Yukawa couplings

→ free parameters (>20)

Renormalisation → free parameters
Traditional way of reducing the number of parameters

SYMMETRY

Celebrated example: GUTs

\[ \text{testable } \sin^2 \theta_W \]

\[ \text{successful } \frac{m_t}{m_b} \]

\[ \text{e.g. minimal } SU(5) \]

However more SYMMETRY (e.g. $SO(10)$, $E(6)$, $E(7)$, $E_8$) does not lead necessarily to more predictions for the SM parameters.

Extreme case: Superstring Theory...
On the other hand

**LEP data** \( \rightarrow \) **\( N = 1 \) SU(5)**

**\( N = 1 \) SU(5) \rightarrow **MSSM**

MSSM best candidate for

**Physics Beyond SM**

But with \( > 100 \) free parameters mostly in its SSB sector.

- Cures problem of quadratic divergencies of the SM (hierarchy problem)
- Restricts the Higgs sector leading to approximate prediction of the Higgs mass
SM with two-Higgs doublets

\[
V = m_1^2 |H_1|^2 + m_2^2 |H_2|^2 + (m_3^2 H_1 H_2^* + h.c.) \\
+ \frac{1}{2} \lambda_1 (H_1^+ H_1)^2 + \frac{1}{2} \lambda_2 (H_2^+ H_2)^2 \\
+ \lambda_3 (H_1^+ H_1)(H_2^+ H_2) + \lambda_4 (H_1 H_2)(H_1^+ H_2^*) \\
+ \left\{ \frac{1}{2} \lambda_5 (H_1 H_2)^2 + \left[ \lambda_6 (H_1^+ H_1) + \lambda_7 (H_1^+ H_2^*) \right] (H_1 H_2) + h.c. \right\}
\]

Supersymmetry provides tree level relations among couplings

\[
\lambda_1 = \lambda_2 = \frac{1}{4} (g^2 + g'^2) \\
\lambda_3 = \frac{1}{4} (g^2 - g'^2), \quad \lambda_4 = -\frac{1}{4} g^2 \\
\lambda_5 = \lambda_6 = \lambda_7 = 0
\]

With \( u_1 = \langle \text{Re } H_1^0 \rangle \), \( u_2 = \langle \text{Re } H_2 \rangle \)

and \( u_1^2 + u_2^2 = (246 \text{ GeV})^2 \), \( \frac{u_2}{u_1} = \tan \beta \)

\( \Rightarrow h^0, H^0, H^\pm, A^0 \)
At tree level

\[ M_{h^0, H^0}^2 = \frac{1}{2} \left\{ M_A^2 + M_Z^2 + \left[ (M_A^2 + M_Z^2)^2 - 4 M_A^2 M_Z^2 \cos^2 2\beta \right] \right\} \frac{1}{2} \]

\[ M_{H^\pm}^2 = M_w^2 + M_A^2 \]

\[ \begin{cases} 
M_{h^0} < M_Z | \cos 2\beta | \\
M_{H^0} > M_Z \\
M_{H^\pm} > M_w 
\end{cases} \]

Radiative corrections

\[ M_{h^0}^2 = M_Z^2 \cos^2 2\beta + \frac{38 \tan^4 \beta}{16 \pi^2 M_w^2} \log \frac{M_{t_1} M_{t_2}}{m_t^2} \]
Grand Unified Theories

The $SU(5)$ model

From group theory (e.g. Slansky, Phys. Repts)

$SU(5) \supset SU(3) \times SU(2) \times U(1)$

$5 = (3,1)_{-2/3} + (1,2)_1$ - fundamental rep $\psi_i$

$10 = (3,2)_{1/3} + (\bar{3},1)_{-4/3} + (1,1)_2$

$24 = (8,1)_0 + (3,2)_{-5/3} + (\bar{3},2)_{5/3}$

$\text{adjoint rep}$

in addition $\psi_{ij} = -\psi_{ji}$

$5 \times 5 = 10 + 15$

$10 \times 10 = \bar{5} + 45 + 50$

$\bar{5} \times 10 = 5 + 45$

$5 \times \bar{5} = 1 + 24$
Recall

Dirac equation in electromagnetic field

\[
(i \gamma \mu \partial_\mu - g \gamma \mu A_\mu(x) - m) \psi(x) = 0
\]

\[
\begin{cases}
(i \gamma \mu \partial_\mu + e \gamma \mu A_\mu(x) - m) \psi^e(x) = 0 & \text{electron} \\
(i \gamma \mu \partial_\mu - e \gamma \mu A_\mu(x) - m) \psi^c(x) = 0 & \text{positron (antiparticle)}
\end{cases}
\]

\[
\rightarrow \quad \psi^c(x) = i \gamma^2 \psi^c(x)
\]

\[
\Rightarrow \quad (\psi_R)^c = (\psi^c)_L = \psi^c_L
\]

Then we can write e.g. the quarks of the first family using the \(SU(3)_c \times SU(2)_L \times U(1)\) quantum numbers and only left-handed 2-comp. fields

\[
\begin{align*}
u_L, d_L &: (3, 2)_1^{\frac{1}{3}} \\
u^c_L &: (\bar{3}, 1)_{-\frac{1}{3}} \\
d^c_L &: (\bar{3}, 1)_{\frac{2}{3}}
\end{align*}
\]
A comparison shows that one family of fermions of the SM can be accommodated in two $SU(5)$ irreps (or three if $\nu_2$ exists):

$\bar{5} : (\psi^i)_L = (d^c_1 d^c_2 d^c_3 \, e^- - \nu_e)_L$

or

$5 : (\psi^i)_R = (d_1 d_2 d_3 \, e^+ - \nu_e^c)_L$

and

$$
\begin{pmatrix}
0 & u^{c^3} - u^{c^2} u_1 & d_1 \\
0 & u^{c^1} u_2 & d_2 \\
0 & u_3 & d_3 \\
0 & 0 & e^+ \\
0 & 0 & 0
\end{pmatrix}
$$

The combination $\bar{5}$ and $10$ is anomaly free.
$SU(5)$ generators

$\{ \lambda^a \}, \ a = 1, \ldots, 24$

is a set of 24 ($5^2 - 1$) unitary, traceless $5 \times 5$ (fundamental rep) matrices with normalisation

$$\text{tr}(\lambda^a \lambda^b) = 2 \delta^{ab}$$

satisfying the CRs

$$\left[ \frac{\lambda^a}{2}, \frac{\lambda^b}{2} \right] = i C^{abc} \frac{\lambda^c}{2} , \ C^{abc} - \text{structure constants}$$

i.e. generalized Gell-Mann matrices

**Examples**

$$\lambda^a = \begin{bmatrix} \lambda^a & 0 & 0 & 0 \\ \lambda^a & 0 & 0 & 0 \\ 0 & 0 & 0 & 0 \\ 0 & 0 & 0 & 0 \end{bmatrix}$$

$a = 1, \ldots, 8$

Gell-Mann generators

(tie corresponding
gauge bosons are the
gluons)
\[ \lambda^{9,10} = \begin{bmatrix} 0 & 0 & 0 & 0 & 0 & 0 \\ 0 & 0 & 0 & 0 & 0 & 0 \\ 0 & 0 & 0 & 0 & 0 & 0 \\ 0 & 0 & 0 & 0 & 0 & 0 \\ 0 & 0 & 0 & 0 & 0 & 0 \\ 0 & 0 & 0 & 0 & 0 & 61,2 \end{bmatrix}, \quad \lambda^{11} = \text{diag} \begin{bmatrix} 0, 0, 0, 1, -1 \end{bmatrix} \]

generators of $SU(3)$
(corresponding the $W^\pm, W_3$)

\[ \lambda^{12} = \frac{1}{\sqrt{15}} \text{diag} \begin{bmatrix} -2, -2, -2, 3, 3 \end{bmatrix} \]
(corresponds the B)

\[ \lambda^{13} = \begin{bmatrix} 0 & 1 & 0 \\ 1 & 0 & 0 \\ 0 & 0 & 0 \end{bmatrix}, \quad \lambda^{14} = \begin{bmatrix} 0 & i & 0 \\ -i & 0 & 0 \\ 0 & 0 & 0 \end{bmatrix}, \text{etc.} \]
Charge Quantization

Property of simple non-abelian groups that the eigenvalues of the generators are discrete (recall e.g. \( SO(3) \))

- In \( SU(5) \) the \( Q \) is one (linear combination) of generators and therefore quantized.
- Since electric charge is additive qu. no., \( Q \) must be some linear combination of the 4 diagonal generators of \( SU(5) \) (rank 4).
- Since \( Q \) commutes with \( SU(3)_c \) generators (2 diagonal; rank 2).
\[ Q = T_3 + \frac{Y}{2} = \frac{1}{2} \lambda'' + \frac{c}{2} \lambda''^2 \]

\( c \) is determined by comparing the eigenvalues of \( \lambda'' \) with the hypercharge \( Y \) values of particles in 5

\[ \Rightarrow \quad c = \left( \frac{5}{3} \right)^{1/2} \]

Then the traceless condition,

\[ \text{Tr} Q = 0 \quad \Rightarrow \quad 3q_d + q_e = 0 \]
(More) SU(5) Physics

- 1st generation of fermions

\[ SU(5) \rightarrow SU(3)_c \times SU(2) \times U(1) \]
\[ 5 = (\bar{3}, 1)_{-\frac{2}{3}} + (1, 2)_{-1} \]
\[ = d^c_L + (v_e, e^-)_L \]
\[ 10 = (3, 2)_{\frac{1}{3}} + (\bar{3}, 1)_{-\frac{4}{3}} + (1, 1)_{2} \]
\[ = (u, d)_L + u^c_L + e^+_L \]

- Gauge bosons

\[ 24 = (8, 1)_{0} + (3, 2)_{-\frac{5}{3}} + (\bar{3}, 2)_{\frac{5}{3}} + (1, 3)_{0} + (1, 1)_{0} \]

\[ \begin{array}{c}
\text{gluons} \\
G \\
\begin{bmatrix}
X_1 Y_1 \\
X_2 Y_2 \\
X_3 Y_3 \\
X^1 X^2 X^3 \\
Y^1 Y^2 Y^3 \\
W_3^{1/2} \\
- W_3^{1/2}
\end{bmatrix}
\end{array} + \frac{B}{160} \begin{bmatrix}
-2 \\
-2 \\
3
\end{bmatrix} \]

\[ \begin{bmatrix}
-2 \\
-2 \\
3
\end{bmatrix} \]
from covariant derivatives

\[ \text{Lint} = - \frac{g_5}{2} \sum_{a} G_{\mu}^{a} \left( \bar{u} \gamma_{\mu} j_{\mu} u + \bar{d} \gamma_{\mu} j_{\mu} d \right) \]

\[ - \frac{g_5}{2} W^{\mu} \left( \bar{Q} \gamma^{\mu} T^{i} Q + \bar{L} \gamma^{\mu} T^{i} L \right) \]

\[ - \frac{g_5}{2} \left( \frac{3}{5} \right)^{1/2} B_{\mu} \sum_{\text{all fermions}} \bar{f} \gamma^{\mu} Y f \]

+ interactions of $X, Y_5$

Note that $g_{\text{strong}} = g_{\text{SU(3)}} = g_5$

however $g_5 X^{1/2} A^{1/2}_\mu = g' Y_B \mu$

and $Y = (\frac{5}{3})^{1/2} X^{1/2}$

$\Rightarrow g' = (\frac{3}{5})^{1/2} g_5$

$\Rightarrow \sin^2 \theta_W = \frac{g'^2}{g_{\text{SU(3)}} + g'^2} = \frac{3}{8}$
Spontaneous Symmetry Breaking

The first symmetry breaking $SU(5) \rightarrow SU(3)_c \times SU(2)_L \times U(1)$ is achieved by introducing a 24-plet of scalars $\Phi(x)$, with potential

$$V(\Phi) = -m_i^2 (\text{tr} \Phi^2) + \lambda_1 (\text{tr} \Phi^2)^2 + \lambda_2 (\text{tr} \Phi^4)$$

for $\lambda_1 > -7/30$, $\lambda_2 > 0$

$$\langle \Phi \rangle = V \begin{pmatrix} 1 & 1 & 1 & -\frac{3}{2} & -\frac{3}{2} \\ 1 & 1 & -\frac{3}{2} & -\frac{3}{2} & -\frac{3}{2} \\ \end{pmatrix}$$

L.-F. Li

where $V^2 = m_i^2 / (15 \lambda_1 + 7/2 \lambda_2)$

The gauge bosons $X, Y$ obtain mass

$$m_X^2 = m_Y^2 = \frac{25}{8} g_5^2 V^2$$
Examining further the interaction of $X, Y$ we find

$$\mathcal{L}_{\text{int}} X Y = -\frac{g_5}{2} \left[ X \bar{\mu} \gamma \left( \frac{d_R \gamma_{\mu} e^c}{e^c} + \frac{d_L \gamma_{\mu} e^c}{e^c} \right) + e \bar{\mu} X \gamma_{\mu} e^c + h.c. \right]$$

$$-\frac{g_5}{2} \left[ Y \bar{\mu} \gamma \left( \frac{d_R \gamma_{\mu} v_R}{v_R} + \frac{d_L \gamma_{\mu} v_R}{v_R} \right) + e \bar{\mu} X \gamma_{\mu} v^c + h.c. \right]$$

leading to proton decay via diagrams such as

$$P \left\{ \begin{array}{c}
\begin{array}{c}
\text{d} \rightarrow \bar{e}^+ \\
\mu \rightarrow \bar{e}^+ \\
\end{array}
\end{array} \right\} \pi^0$$

with strength $\sim \frac{g_5^2}{m_X^2, Y}$

$$\Rightarrow T_p = \frac{m_X^4}{g_5^2 m_P^5}$$

Limits $T_p \geq 10^{31}$, $\Rightarrow m_X, Y \geq 10^{15}$ GeV
Gauge Hierarchy Problems

The second breaking in SU(5), i.e.
$SU(3)_c \times SU(2)_L \times U(1) \rightarrow SU(3)_c \times U(1)_{em}$
is due to a 5-plet of scalars.

Then the complete potential is

$V = V(\Phi) + V(H) + V(\Phi, H)$

5-plet

with

$V(H) = -m_2^2 H^+ H + \lambda_3 (H^+ H)^2$

$V(\Phi, H) = \alpha (H^+ H)(\Phi - \Phi^2) + \beta H^+ \Phi^2 H$

In turn the vev of $\Phi$ changes a bit

$$\langle \Phi \rangle = V \begin{pmatrix} 1 & 1 \\ 1 & -\frac{3}{2} - \xi \\ 1 & -\frac{3}{2} + \xi \end{pmatrix}$$
\[ \langle H \rangle = \frac{v}{\sqrt{2}} \begin{pmatrix} 0 \\ 0 \\ 0 \\ 0 \\ 1 \end{pmatrix} \]

for an appropriate range of parameters.

In turn the gauge boson masses are superheavy $X, Y$: $m_X^2 = m_Y^2 = \frac{25 g_s^2 V^2}{\theta}$

$W, Z$: $m_W^2 \approx \frac{g^2 v^2 (1 + \epsilon)}{2}$, $m_Z^2 \approx \frac{g^2 v^2}{4 \cos^2 \theta_W}$

- to keep the two scales in the theory \( \Rightarrow \epsilon \sim 10^{-28} \)
- the triplet in the H 5-plet can mediate proton decay. Therefore should be superheavy, while the doublet has electroweak scale mass.
Use of Renormalization Group Equations

Our picture is that at scales above $V$ we have a gauge invariant $SU(5)$ theory, which at $\sim V$ breaks down spontaneously to the $SM$ $SU(3)_c \times SU(2)_L \times U(1)$

Then the evolution of the three gauge coupling $g_3, g_2, g_1$, is controlled by the corresponding $\beta$-functions

$$\frac{d g_i}{dt} = \beta_i (g_i), \quad i = 3, 2, 1$$

with $t = \ln \frac{1}{\Lambda M}$
Specifically

\[ b_3 (g_3) = - \frac{g_3^2}{16 \pi^2} \left( 11 - \frac{2}{3} N_f - \frac{N_4}{6} \right) \]

\( N_f \) - number of 4-component colour triplet fermions

\( N_4 \) - number of colour triplet Higgs bosons

\[ b_2 (g_2) = - \frac{g_2^2}{16 \pi^2} \left( \frac{22}{3} - \frac{2}{3} N_f - \frac{N_4}{6} \right) \]

\[ b_1 (g_1) = + \frac{2}{3} N_f \frac{g_1^2}{16 \pi^2} + \text{Higgs boson contr.} \]

In addition we have the boundary condition: 

\[ g_3 (M_X) = g_2 (M_X) = g_1 (M_X) = g' (1.5) = g_5 \]

We find

\[ M_X = 2.1 \times 10^{14} \times (1.5)^{1/3} \left( \frac{\Lambda_{MS}}{0.16 \text{ GeV}} \right) \]

\[ \Rightarrow \sin^2 \theta_W (M_W) = 0.214 \pm 0.003 \pm 0.006 \ln \left( \frac{0.16 \text{ GeV}}{\Lambda_{MS}} \right) \]

Experiment \( \sin^2 \theta_W = 0.23149 \pm 0.00017 \)
Fermion masses in $SU(5)$

Introduce Yukawa couplings

$f_1 10 + f_2 10f \bar{5} \bar{5} \bar{4}$

$\langle 5_H \rangle = (0, 0, 0, 0, 0, \nu/\sqrt{2})$

$\Rightarrow \frac{\nu}{\sqrt{2}} f_1 \bar{u} u + \frac{\nu}{\sqrt{2}} f_2 (d \bar{d} + e \bar{e})$

i.e. \[ m_e = \mu \]

holding at mass scales that $SU(5)$ is a good symmetry, subject to significant renormalization corrections

\[ m_b / m_t \approx 3 \] at $\mu = \mu_{\text{th}} \approx 10 \text{ GeV}$

but \[ m_{\mu} / m_e = m_{\text{3rd}} / m_{\text{1st}} \]

which can be improved to \[ m_{\mu} / m_e = 9 \]

by introducing a 45-plet of scalars
Supersymmetric $SU(5)$

- Motivation: Try to solve the hierarchy problem. Recall that tree level parameters were tuned to an accuracy $10^{-26}$ to generate $M_X/m_W \sim 10^{12}$ in the scalar potential of $SU(5)$. This relation is destroyed by renormalization effects. It has to be enforced order by order in pert. th.

  Supersymmetry can solve the technical problem. If it is exact, the mass parameters of the potential (in fact the whole superpotential) do not get renormalized. (non-renorm. th.)
When susy is broken corrections are finite and calculable (\(a \Delta m^2\) square mass splitting in supermultiplet).

- RGEs estimation of the GUT scale and proton life-time.
- Gauge bosons the same
- Larger number of spinors and scalars
  - Smaller, in absolute value, \(b\)-function and therefore slower variation of the asymptotically free couplings.

\[
M_x, \text{susy} \lesssim 10^{16} - 10^{17} \text{ GeV}
\]

\[
\sin^2 \Theta_W(m_W) = 0.232
\]

Proton decay is out of reach with usual gauge boson exchange.
There exist new diagrams (based on dimension 5 operators, dressed by gauginos) that become dominant. The resulting proton lifetime is compatible with presently known bounds.

An important outcome is that the dominant decay mode is different from the $SU(5)$ model and is

$$P \rightarrow K^+ \nu$$

(due to the anti-symmetry in the colour index that appears, the resulting operators are flavour non-diagonal).
Quantum
Reduction of Couplings

Consider a GUT with
\( g \) - gauge coupling
\( g_i \) - other couplings (Yukawas, self-couplings)

Any relation among the couplings
\( \Phi(g, g, \ldots) = \text{const} \)

which is \( RGI \) should satisfy

\[
\frac{d}{dt} \Phi = 0, \quad t = \text{phys}
\]

\[
\frac{d}{dt} \Phi = \frac{\partial \Phi}{\partial g} \frac{dg}{dt} + \sum_i \frac{\partial \Phi}{\partial g_i} \frac{dg_i}{dt} = 0
\]

which is equivalent to

\[
\frac{dg}{g} = \frac{dg_1}{g_1} = \frac{dg_2}{g_2} = \ldots \quad \text{characteristic system}
\]
Demand power series solution to the REs

\[ S_c = \sum_{n=0}^{\infty} \rho^{n+1} g^{2n+1} \]

Remarkably, uniqueness of these solutions can be decided already at 1-loop!

Assume

\[ \bar{b}_i = \frac{1}{16 \pi^2} \left[ \sum_{j,k,l} \bar{b}_i^{(1)} j k l g g g + \sum_{j \neq l} \sum_{i} \bar{b}_i^{(1)} j i j + \cdots \right] \]

\[ \bar{b}_g = \frac{1}{16 \pi^2} \bar{b}_g^{(1)} g^3 + \cdots \]

Assume \( \rho_i^{(n)} \), \( n \leq r \) have been uniquely determined

To obtain \( \rho_i^{(r+1)} \), insert \( g_i \) in REs and collect terms of \( O(g^{2r+1}) \)
\[ \sum_{l \neq g} M(r)_i^l \rho_{(r+1)}^l = \text{lower order quantities} \]

known by assumption

where

\[ M(r)_i^l = 3 \sum_{j, k \neq g} \delta_{(2)}{i}^{j, k} p_j^{(1)} p_k^{(1)} \rho_{(1)}^l - (2r+1) \delta_{(1)}{i}^{l} \delta_{(g)}{i} \]

\[ 0 = \sum_{j, k, l \neq g} \delta_{(1)}{i}^{j, k} p_j^{(1)} p_k^{(1)} \rho_{(1)}^l + \sum_{l \neq g} \delta_{(1)}{i}^{l} \rho_{(1)}^l - \delta_{(g)}{i} \rho_{(1)}^l \]

\[ \Rightarrow \text{for a given set of } p_{i}^{(1)}, \text{ the } p_{i}^{(n)} \text{ for all } n > 1 \text{ can be uniquely determined if} \]

\[ \det M(n)_i^l \neq 0 \]

for all \( n \)
Consider an $SU(N)$ (non-susy) theory with

$\phi^i(N), \bar{\phi}^i(N)$ - complex scalars

$\psi^i(N), \bar{\psi}^i(N)$ - Weyl spinors

$\gamma_a (a = 1, \ldots, N^2 - 1)$

\[
\mathcal{L} = -\frac{1}{4} F_{\mu\nu} F^{\mu\nu} + i \sqrt{2} \left[ g_Y \bar{\psi}^i \gamma^a T^a \phi^i + \bar{\psi}^i \gamma^a T^a \phi + h.c. - V(\phi, \bar{\phi}) \right]
\]

\[
V(\phi, \bar{\phi}) = \frac{1}{4} J_1 (\phi^i \phi^*_i)^2 + \frac{1}{4} J_2 (\phi^i \bar{\phi}^*_i)^2 + J_3 (\phi^i \phi^*_i) (\phi^j \bar{\phi}^*_j) + J_4 (\phi^i \bar{\phi}^*_j) (\phi^j \bar{\phi}^*_i)
\]

Searching for power series solution of the R.E.s we find

$g_Y = \bar{g}_Y = g; J_1 = J_2 = \frac{N-1}{n} g^2; J_3 = \frac{1}{2n} g^2; \lambda_4 = -\frac{1}{2} g^2$

i.e. SUSY
$N = 1$ gauge theories

Consider a chiral, anomaly free $N = 1$ globally supersymmetric gauge theory based on a group $G$ with gauge coupling $g$.

Superpotential

$$W = \frac{1}{2} \epsilon_{ij} \phi^i \phi^j + \frac{1}{6} C_{ijk} \phi^i \phi^j \phi^k$$

$m_{ij}, C_{ijk}$ - gauge invariant tensors
$q^i$ - matter fields transforming as an irrep $R_i$ of $G$. 

Renormalization constants associated with \( W \)

\[
\phi^0_i = (Z^j_i)^{1/2} \phi^j, \quad m^0_{ij} = Z^{ij}_{i'j'} m_{i'j'}, \quad C_{ijk} = Z^{ijk}_{i'j'k'} C_{ij'k'}
\]

\( N = 1 \) non-renormalization thus ensures absence of mass and cubic-in-field infinities

\[
Z^{i j k}_{i' j' k'} Z^{j' k l}_{i'' j'' k''} Z^{k l m}_{i''' j''' k'''} = \delta^{i'}_{i} \delta^{j'}_{j} \delta^{k'}_{k},
\]

\[
Z^{i j i'} Z^{j' k l} Z^{k l m}_{i'' j'' k''} = \delta^{i'}_{i} \delta^{j'}_{j},
\]

(1 in the background field method)

\[
Z^g Z^{i j}_{i' j'} = 1
\]

Only surviving infinities are \( Z^g_{i j} \)\( (Z^i_j) \)

i.e. one infinity for each field.
The 1-loop β-function of the gauge coupling is

\[ β_g^{(1)} = \frac{dg}{dt} = \frac{g^3}{16\pi^2} \left[ \sum_{i} C(R_i) - 3C_2(G) \right] \]

\( C(R_i) \) - Dynkin index of \( R_i \)
\( C_2(G) \) - quadratic Casimir of the adjoint rep.

β-functions of \( C_{ijk} \), by virtue of the non-renormalization theorem, are related with the anomalous dim. matrix \( \delta_{ij} \) of \( g^a \):

\[ b_{ijk} = \frac{dC_{ijk}}{dt} = C_{ije} \delta_k^e + C_{ike} \delta_j^e + C_{ihek} \delta_j^e \]

\[ \gamma^{(1)}_{ij} = z^{-1/2}_{ik} \frac{d}{dt} z^{1/2}_{ij} \]

\[ = \frac{1}{32\pi^2} \left[ C_{ike} C_{ike} - 2 g^2 C_{2(R_i)} \delta_{ij} \right] \]

\( C_{2(R_i)} \) - quadratic Casimir of \( R_i \)
\( C_{ijk} = C^*_{ijk} \)
\[ g_4^{(3)} = \frac{1}{(16\pi^2)^2} \sum_i \left[ \sum_k \left( l(R_i) - 3C_2(G) \right) \right] \]

\[ - \frac{1}{(16\pi^2)^2} \frac{g^3}{r} \sum_i C_2(R_i) \left[ C_{ik}^l \left( \frac{C_{klm}^i}{r} - 2g^2C_2(R_i) \right) \right] \]

\[ \gamma^{(2)i} = \frac{1}{(16\pi^2)^2} \left[ \sum C_2(R_i) \left( \sum l(R_i) - 3C_2(G) \right) \right] \]

\[ - \frac{1}{(16\pi^2)^2} \frac{1}{2} \left[ C_{ik}^l \left( C_{jkm} + 2g^2(R)_m^i \right) \right] \]

\[ \left[ C_{mpq} \left( C_{epq} - 2\delta^e_\nu g^2 C_2(R_i) \right) \right] \]

\[ g_{NSVZ}^{(5)} = \frac{g^3}{16\pi^2} \left[ \frac{\sum C(R_i) (1 - 2\delta_i) - 3C_2(G)}{1 - g^2C_2(G) / 8\pi^2} \right] \]

Novikov - Shifman - Vainshtein - Zakharov
Finite Unification

Old days...
...divergences are "hidden under the carpet" (Dirac, Lects on Q.F.T, '64)

Recent years...
...divergences reflect existence of a higher scale where new degrees of freedom are excited. Not just artifacts of pert. th.

However the presence of quadratic divergences means that physics at one scale are very sensitive to unknown physics at higher scales.
SUSY theories which are free of quadratic divergences in spite of any experimental evidence...

Natural to expect that beyond unification scale the theory should be completely finite.

- $N = 4 \Rightarrow$ finite to all orders in pert.
- $N = 2 \Rightarrow$ only 1-loop contributions to $\beta$-function. Possible to arrange the spectrum so that theory is finite.
Multiplicities for massless irreducible reps with maximal helicity 1

<table>
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<tr>
<th>N</th>
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\[ N=2 : b(g) = \frac{2g^3}{(4\pi)^2} \left( \sum_i T(R_i) - C_2(G) \right) \]

e.g. \( SU(N) \) with \( 2N \) fundamental reps \( \rightarrow b(g) = 0 \)

\( SU(5) : p(\bar{5}+\bar{5}) ; q(\bar{10}+\bar{10}) ; r(15+\bar{15}) \)

with \( p + 3q + 7r = 10 \)

\( SO(10) : p(10+\bar{10}) ; q(16+\bar{16}) \)

with \( p + 2q = 8 \)

\( E_6 \) : 4(27+\bar{27}) \)
Finite Unified Theories

$N = 1$

- 1-loop finiteness conditions
  \[ \eta^{(1)}_{ij} = 0 \]
  \[ \gamma^{(1)}_{ij} = 0 \] - anomalous dimensions of all chiral superfields

- Exists complete classification of all chiral $N=1$ models with
  \[ \eta^{(1)}_{ij} = 0 \] Hamidi - Patera - Schwarz
  Jiang - Zhou

- 1-loop finiteness

- 2-loop finiteness
  Parkes - West
  Jones
  Mezincescu
Exist simple criteria that guarantee all loop finiteness (vanishing of all-loop beta functions)

All-loop finite \text{SU}(5) \implies \text{top quark mass} \sqrt{2}

Susy sector

1-loop finiteness cords \text{SUSY}

(requires in particular universal soft susy scalar masses

\begin{equation}
(m^2)_j = \frac{1}{3} M M^* \delta_{ij}
\end{equation}
1-loop finiteness $\Rightarrow$ 2-loop finiteness

Reduction of couplings

- Extension of method in SSB sector
- Application in min susy SU(5)

1-loop sum rule for soft scalar masses in non-finite susy ths.

2-loop sum rule for soft scalar masses in finite ths.

* All-loop RG I relations in finite and non-finite ths.
All-loop sum rule for soft scalar masses is finite and non-finite ths.

** SU(5) FUTs **

- Prediction of s-spectrum in terms of few parameters starting from several hundreds GeV.

- The LSP is neutralino (see e.g., Kazakov et. al., Yoshioka)

- Radiative E-W breaking (see e.g., Brignole, Banerji, Mann)

- No funny colour, charge (see e.g., Casas et. al.)

- Prediction of Higgs masses

  Lightest ~ 118 - 129 GeV

Similar results also for win SUSY SU(5)
Consider a chiral, anomaly free, $N=1$ gauge theory with group $G$. The superpotential is
\[
W = \frac{1}{6} \, Y^{i j k} \, \Phi_i \Phi_j \Phi_k + \frac{1}{2} \, \mu^{i j} \Phi_i \Phi_j
\]

$Y^{i j k}$: gauge invariant
$\mu^{i j}$: Yukawa couplings

$\Phi_i$: matter superfields in irreducible reps of $G$

Necessary and sufficient conditions for $N=1$ 1-loop finiteness

- Vanishing of $\delta g^{(1)}$ implies

\[
\sum_i \ell(R_i) = 3 \, C_2(G)
\]

$\ell(R_i)$: Dynkin index of $R_i$
$C_2(G)$: Quadratic Casimir of $G$ (adjoint)

$\Rightarrow$ Selection of the field content (representations) of the theory
Vanishing of $\chi^{(i)}_{j}$ implies

$$Y_{\mu}^{\mu} \ Y_{\nu}^{\nu} = \sum_{i} \delta_{\mu i} \ g^{2} C_{2}(R_{i})$$

Yukawa and gauge couplings are related.

Note: $\mu$'s are not restricted.

Appearance of $U(1)$ is incompatible with 1st cond.

2nd cond. forbids the presence of singlets with nonvanishing couplings.

$\therefore \Rightarrow$ SUSY by G-invariant soft terms.
1-loop finiteness conditions necessary and sufficient to guarantee 2-loop finiteness.

1-loop finiteness conditions ensure that \( g^{(3)}_\gamma \) in 3-loops vanishes but in general \( \gamma^{(3)} \) does not.

Grisaru - Milewski - Zanon 
Part Ies - West

What happens in higher loops?

So far 1-loop finiteness conditions (or \( \gamma \)s) are telling us

\[
\gamma_{ijk} = \gamma_{ijk}(g) \\
\delta_{ijkl} = 0
\]
Necessary and sufficient conditions for vanishing $B_j$ and $B_{ijk}$ to all orders

1. $B_j^{(1)} = 0$

2. $Y_j^{(1)} = 0$

3. $B_{ijk} = B_j \frac{dY_{i,jk}}{dg}$

admit power series solution which in lowest order is a solution of order 2.

3'. There exist solutions to $Y_j^{(1)} = 0$
of the form

$Y_{i,jk} = p_{ijk} g$, $p_{ijk}$ - complex

4. These solutions are isolated and non-degenerate considered as solutions of $B_Y = 0$
Recall
R-invariance, axial anomaly

In massless $N=1$ the $U(1)$ chiral transformation $Q$: $A \mu \rightarrow A \mu$, $j \rightarrow e^{-i\alpha j}$, $\phi \rightarrow e^{-i\frac{2}{3}\alpha}\phi$, $\chi \rightarrow e^{i\frac{1}{3}\alpha}\chi$, ...

\[\psi D = \left(\begin{array}{c} \frac{i}{\chi} \\
\end{array}\right) \rightarrow e^{i\alpha \gamma_5}\psi_D\]

Noether current $J\mu = \bar{\psi} \gamma^\mu \gamma^5 \psi + \ldots$

$\rightarrow \partial \mu J^\mu = r(\epsilon^{\mu
\nu\rho\sigma} F_{\nu\rho} F_{\sigma\mu} + \ldots)$

$r = G^{(2)}_{1}$!

Only 1-loop contributions due to non-renormalization thus.

Adler, Bardeen, Jackiw, Pi, Shei, Zee
Supercurrent
\[ J^I = \{ J_R^I, Q^I, T^I \} \]

vector super multiplet

associated to R-invariance

Ferrara + Zumino

supercurrent is represented as vector superfield
\[ V_\mu(x, \theta, \bar{\theta}) = \tilde{Q}_\mu(x) - i \tilde{\theta}^\alpha Q^{\alpha \mu}(x) + i \bar{\theta}^\dot{\alpha} \bar{Q}^{\dot{\alpha} \mu}(x) - 2(\theta \delta^\nu \bar{\theta}) T_{\mu \nu}(x) + \cdots \]

\[ J_R^J = J_R^J \]

\[ J_R^J = J_R^J + 0(\delta) \]

In addition
\[ S = \{ g^F_\mu \{ F_{\mu \nu} + \cdots, \delta g \epsilon_{\mu \nu \rho \sigma} F_{\rho \sigma} F_{\mu \nu} + \cdots \} \]

super trace anomaly

trace anomaly of SUSY current

trace anomaly of SUSY current
There is a relation, whose structure is independent from the renormalization scheme, although individual coefficients (except the 1-loop values of $b$-functions) may be scheme dependent.

\[ r = 6g (1 + x_g) + \delta_{ijk} x^{ijk} - \delta_A r_A \]

Linear combinations of anomalous divergences

Thm: If
(i) no gauge anomaly
(ii) \( B^{(2)} = 0 \) i.e. no $B$-current anomaly
(iii) \( \gamma^{(1)} i = 0 \) implies also \( r_A = 0 \)
(iv) exist solutions to \( \gamma^{(1)} = 0 \) of the form \( C_{ijk} = P_{ijk} + g, P_{ijk} \)-complex
(v) these solutions are isolated, non-degenerate
Then each of all solutions can be uniquely extended to a formal power series in $g$, and the $N = 1$ YM models depend on the single coupling constant $g$ with a $g^k$-function vanishing to all orders.

Proof: Inserting $g_{ijk} = \frac{\partial g_{ijk}}{\partial g}$ in the identity and taking into account the vanishing of $r, r^k \rightarrow 0 = g^k (1 + O(\frac{1}{r}))$ its solution (as formal power series in $\frac{1}{r}$) is: $g = 0$ and $\mathcal{B}_{ijk} = 0$ too.
2-loop RGEs for SSB parameters

Consider $N=1$ gauge theory with

$$W = \frac{1}{6} Y^{ijk} \phi_i \phi_j \phi_k + \frac{1}{2} \mu_{ij} \phi_i \phi_j$$

and SSB terms

$$-L_{SB} = \frac{1}{6} h^{ijk} \phi_i \phi_j \phi_k + \frac{1}{2} b^{ij} \phi_i \phi_j$$

$$+ \frac{1}{2} (m_2)^2 \phi_i^* \phi_i + \frac{1}{2} M \phi_i^* \phi_i + h.c.$$ 

• 1-loop finiteness conditions

$$h^{ijk} = - M Y^{ijk}$$

$$(m_2)^i_j = \frac{1}{3} M M^* \delta^i_j$$

in addition to $\theta^{(1)}_i = \gamma^{(1)}_i = 0$

• 1-loop finiteness

$\rightarrow$ 2-loop finiteness
Assuming
• $\theta_i^{(1)} = y_i^{(1)} s_i = 0$
• the reduction eq
  $\xi_{ijk} = g \sum \frac{dY_{ijk}}{dg}$
  admits power series solution
  $Y_{ijk} = g \sum_{n=0}^{\infty} P_{ijk}^n g^n$
• $(m^2)_j^i = m^2 \delta_j^i$

\[ \left( m_{ij}^2 + m_j^2 + m_k^2 \right)/M^* = 1 + \frac{g^2}{16\pi^2} \Delta^{(2)} \]

for $i, j, k$ with $P_{ijk}^{(1)} \neq 0$

where $\Delta^{(2)} = \sum_{l} \left( \frac{m_i^2}{m^*} - \frac{1}{3} \right) C(l)$
• $\Delta^{(2)} = 0$ for $N = 4$ with STTr cond
• $\Delta^{(2)} = 0$ for the $N=1$, $SU(5)$ FUTs!
The SU(5) finite model

Kapetanakis, Mondragon, Z
Kobayashi, Kubo, Mondragon, Z

Content

\[ 3(\bar{5} + 10) + 4(5 + \bar{5}) + 24 \]

fermion supermultiplets

\[ H^a \bar{H}^\alpha \]

scalar supermultiplets

Jones-Raby
Hamidi-Schwarz
Accinos et al.
Kazarov
Babu-EnKhabi
Gogoladze

\[ W = \frac{3}{2} \sum_{i=1}^{3} \left[ \frac{1}{3} g^u \Lambda_{101} i_{101} H_i + g^d_{101} i_{101} \bar{5}_i H_i \right] \]

\[ + \frac{1}{3} g^u \Lambda_{23} i_{102} \bar{H}_4 + \frac{1}{3} g^d_{23} i_{102} \bar{5}_3 \bar{H}_4 + \frac{1}{3} g^d_{32} i_{103} \bar{5}_2 \bar{H}_4 \]

\[ + \sum_{\alpha=1}^{3} g^\alpha H_\alpha 24 \bar{H}_\alpha + S^{\gamma/3 (24)^3} \]

(with enhanced discrete symmetry after reduction of couplings)
We find

$$G^{(1)}_0 = 0$$

$$G^{(1)}_i = \frac{1}{16\pi^2} \left[ -\frac{96}{5} g^2 + \sum_{b=1}^{\frac{4}{3}} (g_{ib})^2 + 3 \sum_{j=1}^{\frac{3}{5}} (g_{ja})^2 \\
+ \frac{24}{5} (g_{ix})^2 + 4 \sum_{b=1}^{\frac{4}{3}} (g_{ib})^2 \right] g_{ix}$$

$$G^{(1)}_d = \frac{1}{16\pi^2} \left[ -\frac{84}{5} g^2 + 3 \sum_{b=1}^{\frac{4}{3}} (g_{ib})^2 + \frac{24}{5} (g_{ix})^2 \\
+ 4 \sum_{j=1}^{\frac{3}{5}} (g_{ix})^2 + 6 \sum_{b=1}^{\frac{4}{3}} (g_{ib})^2 \right] g_{ix}$$

$$G^{(1)}_s = \frac{1}{16\pi^2} \left[ -30 g^2 + 6 \sum_{b=1}^{\frac{3}{5}} (g_{ib})^2 + 3 \sum_{a=1}^{\frac{4}{3}} (g_{ia})^2 \right] g_{s}$$

$$G^{(1)}_s = \frac{1}{16\pi^2} \left[ -\frac{98}{5} g^2 + 3 \sum_{i=1}^{\frac{3}{5}} (g_{ia})^2 \\
+ 4 \sum_{c=1}^{\frac{3}{5}} (g_{ic})^2 + \frac{48}{5} (g_{ix})^2 \\
+ \sum_{b=1}^{\frac{4}{3}} (g_{ib})^2 + \frac{21}{5} (g_{ix})^2 \right] g_{s}$$
Considering $g$ as the primary coupling, we solve the Red. Eqs.

$$\delta g = 6a \frac{d\delta}{d\delta a}$$

requiring power series ansatz.

$$\Rightarrow \begin{align*}
(\delta_{ii})^2 &= \frac{8}{5} g^2 + \cdots, \\
(\delta^{-1}_{ii})^2 &= \frac{6}{5} g^2 + \cdots \\
(\delta^1)^2 &= \frac{15}{7} g^2 + \cdots, \\
(\delta^f_4)^2 &= g^2, \\
(\delta^f_\alpha)^2 &= 0 + \cdots (\alpha = 1, 2, 3)
\end{align*}$$

Higher order terms can be uniquely determined.

$$\Rightarrow \text{All 1-loop } B\text{-functions vanish}$$
Moreover

All 1-loop anomalous dimensions of chiral fields vanish.

\[ \Gamma_{10^i} = \frac{1}{16\pi^2} \left[ -\frac{36}{5} \, g^2 + 3 \sum_{b=1}^{4} (\delta_{c\, b})^2 + 2 \sum_{b=1}^{4} (\bar{q}_{ib})^2 \right] \]

\[ \Gamma_{5^i} = \frac{1}{16\pi^2} \left[ -\frac{24}{5} \, g^2 + 4 \sum_{b=1}^{4} (\delta_{c\, b})^2 \right] \]

\[ \Gamma_{H^a} = \frac{1}{16\pi^2} \left[ -\frac{24}{5} \, g^2 + 3 \sum_{i=1}^{3} (\bar{g}_{ia})^2 + \frac{24}{5} (\bar{q}_{ia})^2 \right] \]

\[ \Gamma_{\bar{H}^a} = \frac{1}{16\pi^2} \left[ -\frac{24}{5} \, g^2 + 4 \sum_{i=1}^{3} (\bar{g}_{ia})^2 + \frac{24}{5} (\bar{q}_{ia})^2 \right] \]

\[ \Gamma_{24^i} = \frac{1}{16\pi^2} \left[ -\frac{10}{5} \, g^2 + \sum_{a=1}^{4} (\bar{g}_{ia})^2 + \frac{21}{5} (\bar{q}_{ia})^2 \right] \]

\[ \Rightarrow \text{Necessary and sufficient units for finiteness to all orders are satisfied} \]
• SU(5) breaks down to the standard model due to $\langle 24 \rangle$

• Use the freedom in mass parameters to obtain only a pair of Higgs fields light, acquiring v.e.v.

• Get rid of unwanted triplets rotating the Higgs sector (after a fine tuning)

  see Quirós et al., Kazakov et al.

  Yoshio Ka

• Adding soft terms we can achieve SUSY breaking.
1) Gauge Couplings Unification

\[ \sin^2 \theta_W, \alpha_{em} \rightarrow \alpha_3(M_Z) \]

2) Bottom-Tau Yukawa Unif.

**SU(5)-type**

\[ m_t \sim 100-200 \text{ GeV} \]

Barger et al.

3) Top-Bottom-Tau Yuk Unif.

\[ h_t^2 = \frac{4}{3} h_{b\tau}^2 \quad \text{in SU(5)-FUT} \]

Similar to SO(10)

Ananthanarayan et al.

4) Gauge-Top-Bottom-Tau Unif.

e.g. FUT-SU(5): \[ h_t^2 = \frac{8}{5} g_u^2; \quad h_{b\tau}^2 = \frac{6}{5} g_u^2 \]
<table>
<thead>
<tr>
<th>$M_s$ [GeV]</th>
<th>$\alpha_{3(5f)}(M_Z)$</th>
<th>$\tan \beta$</th>
<th>$M_{GUT}$ [GeV]</th>
<th>$M_b$ [GeV]</th>
<th>$M_t$ [GeV]</th>
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<td>$2.2 \times 10^{16}$</td>
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**FUTA**

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<tr>
<td>$1.2 \times 10^3$</td>
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**FUTB**

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<th>$M_{GUT}$ [GeV]</th>
<th>$M_b$ [GeV]</th>
<th>$M_t$ [GeV]</th>
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</table>

**MIN SU(5)**

The predictions for the three models for different $M_s$

With theoretical corrections and uncertainties $\sim 4\%$

$M_t = 173.8 \pm 5 \text{ GeV}$, $178.0 \pm 4.3 \text{ GeV}$

CDF + D0
Model A

Similar behaviour holds for Model B too

\[ m_\tau^2 \text{ and } m_\chi^2 \text{ for the universal choice of soft scalar masses} \]
Model A

$M_{\text{suby}} = 0.3$ TeV

The empty region yields a neutralino as LSP.
3D) Predictions for the light Higgs boson

- **green**: consistent with $B$ physics constraints
- **red**: agreement with (loose) CDM bound

$$118 \text{ GeV} \leq M_h \leq 129 \text{ GeV} \quad \text{(incl. theor. unc.)}$$

$\Rightarrow$ "easy" to find for LHC (but "only" SM-like ...)

*Sven Heinemeyer, LPT/Orsay particle physics seminar, 26.02.2008*
Typical mass spectrum for FUTB-

<table>
<thead>
<tr>
<th>$m_t$</th>
<th>172</th>
<th>$\bar{m}_b(M_Z)$</th>
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Sven Heinemeyer, LPT/Orsay particle physics seminar, 26.02.2008
<table>
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<td>Mbot(MZ)</td>
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FUTB, $\mu < 0$
Summer School and Workshop on the Standard Model and Beyond
September 8 - 17

TR33 - Summer Institute: Particles and the Universe
September 16 - 22

XVIII European Workshop on String Theory
September 19 - 27

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